



# ATOMIC PHYSICS

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# ATOMIC PHYSICS

BY

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*Authorized Translation from the German Edition by*

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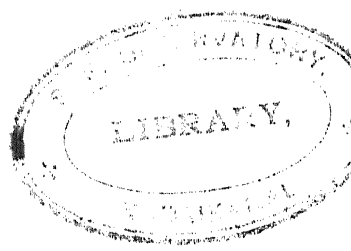
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*Dedicated with gratitude  
to my friend*

ERICH ROSENBERG



## PREFACE

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The German edition of this book, which appeared in 1933 under the title *Moderne Physik*, had its origin in a course of lectures which I gave at the Technical College, Berlin-Charlottenburg, at the instance of the Society of Electrical Engineers, and which were reported by Dr. F. Sauter. The choice of material was not left entirely to myself, the Society making certain suggestions which took into account the previous knowledge and the professional predilections of its members.

A year or two later, the idea of an English edition was mooted. But by this time the title of the book had become inappropriate, for in the interval the growing point of physics had shifted, and interest had become centred in the physics of the nucleus, of which only a bare sketch had been given in the lectures. The change in the title, however, has nothing to do with this, but is due to the fact that the publishers already have a book entitled *Modern Physics* on their list of publications.

I have taken considerable trouble to secure that the book should be brought thoroughly up to date. A new chapter has been introduced, containing an account of the great advances of the last few years—the discovery of new particles, and the elucidation of the properties of nuclei. The original chapters have also been extended by the addition of new sections and paragraphs.

As a theoretical physicist, I have naturally placed the theoretical interpretation of phenomena well in the foreground. The text itself, however, contains only comparatively simple discussions of theory. Proofs, short but complete, of the formulæ employed are collected in a series of Appendices.

In the lectures two years ago, I emphasized with a certain pride the successes of the theory—now I have rather to direct attention to the immense progress which has been made on the experimental side. But I have not been able to resist the temptation to mention a few ideas of my own, which I hope may be of some assistance in stirring the theory into renewed activity.

A new feature of the book is the introduction of dates. Their

purpose is to serve as a historical guide, the years given not being always the first date of publication of an idea or discovery, but that of a specially important paper on the subject. I apologize for not having been able to quote all important papers and authors.

I am much indebted to many helpers. Dr. Cockcroft (Cambridge) and Dr. Nordheim (Haarlem) have read some chapters, which their advice has enabled me to improve. Mr. Pryce, Mr. Weiss and Mr. Goldhaber (Cambridge) have assisted me with the proof-reading. I owe many thanks to Professor P. M. S. Blackett for permission to use some of his beautiful photographs. I am particularly indebted to the translator, Dr. John Dougall, for his painstaking work. He has not only given an adequate rendering of the German, but has also kept a careful look out for misprints or slips in the text and formulæ.

MAX BORN.

CAMBRIDGE, *February, 1935.*

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# ATOMIC PHYSICS

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## CHAPTER I

### Kinetic Theory of Gases

#### 1. Atomic Theory in Chemistry.

In present-day physics, the concepts of energy and matter are connected in the most intimate way with atomic theory. We therefore begin with a brief discussion of the rise of ideas relating to the atom. The source of these ideas, we know, is chemistry. They suggest themselves almost inevitably when we try to interpret the simple regularities which are at once disclosed when the masses of the substances transformed in chemical reactions are determined quantitatively with the balance. It is established in the first place, that in a reaction the total weight remains unchanged. Secondly, it is found that substances combine only in fixed simple proportions by weight, so that a definite weight of one substance can only enter into reaction with definite weights of a second substance; and the ratio of these weights is independent of the external conditions, such as, for example, the proportion by weight in which the two substances may have been mixed. These regularities are expressed in the language of the chemist in the *laws of constant and of multiple proportions* (Proust, 1799; Dalton, 1808). E.g.:

1 gm. hydrogen combines with 8 gm. oxygen to form 9 gm. water,

1 gm. hydrogen combines with 35.5 gm. chlorine to form 36.5 gm. hydrogen chloride.

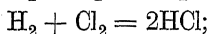
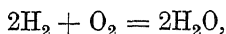
An example of the law of multiple proportions is given by the compounds of nitrogen and oxygen: 7 gm. nitrogen combine with

1 × 4	gm. oxygen to form	11 gm. nitrous oxide,
2 × 4	„ „ „	15 „ nitric oxide,
3 × 4	„ „ „	19 „ nitrous anhydride,
4 × 4	„ „ „	23 „ nitrogen dioxide,
5 × 4	„ „ „	27 „ nitric anhydride.

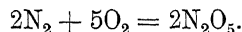
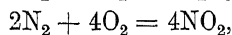
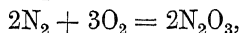
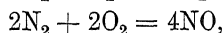
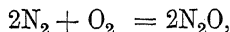
In the case of gases, simple laws hold not only for the weights of the reacting substances but also for their *volumes* (Gay-Lussac, 1808). Thus (at constant pressure)

2 vols. hydrogen combine with 1 vol. oxygen to form 2 vols. water vapour,  
1 vol. hydrogen combines with 1 vol. chlorine to form 2 vols. hydrogen chloride.

The numbers expressing the ratios by volume are those which appear in the corresponding chemical formulæ. In the preceding examples, for instance, we have



and



These facts may be interpreted as follows, as was done by Avogadro: every gas consists of a great number of particles, its atoms or molecules; and equal volumes of all gases, at the same temperature and pressure, contain the same number of molecules.

The significance of this principle in relation to the laws of chemical reactions may be illustrated from the above examples. The fact that two volumes of hydrogen combine with one volume of oxygen to form two volumes of water vapour is (Avogadro, 1811) equivalent to the statement that two molecules of hydrogen combine with one molecule of oxygen to form two molecules of water. Similarly, the combination of one part by weight of hydrogen with eight parts by weight of oxygen to form nine parts by weight of water means that a molecule of oxygen must be eight times, and two molecules of water nine times, as heavy as two molecules of hydrogen.

We are thus led to the concepts of *molecular weight* and *atomic weight*. These are, respectively, the weights of a molecule and of an atom of the substance in question. They are not measured in grammes, but with reference to a standard (ideal) gas, the atomic weight of which is put equal to 1; and it has been agreed to define this, not so that  $\text{H} = 1$ , but so that  $\text{O} = 16$ ; this convention has turned out very lucky (because of the existence of the heavy isotope of hydrogen, p. 59). We shall denote the molecular weight measured in this way by  $\mu$ .

That quantity of a substance, the weight of which is  $\mu$  gm., is called a *mole* (even when the substance is not capable of existence in the chemical sense). A mole of oxygen atoms therefore weighs 16 gm., but a mole of oxygen molecules weighs 32 gm. From this definition of the mole it follows that the quantity "1 mole" always contains the same number of molecules. This number of molecules per mole plays a great part in the kinetic theory of gases; we denote it by  $L$  after the physicist Loschmidt, who first determined it (1865). According to the usual custom with regard to scientific nomenclature it should be called "Loschmidt's number", and this is the expression used in the German literature, but in other languages it is called *Avogadro's number*. Its value is (see p. 22)

$$L = 6.06 \times 10^{23} \text{ mole}^{-1}.$$

As a consequence of Avogadro's law, 1 mole of any gas at a given pressure  $p$  and a given temperature  $T$  always occupies the same volume; for a pressure of 760 mm. of mercury and a temperature of  $0^\circ \text{C}$ . the volume is 22.4 litres.

We add here an explanation of a number of symbols which will be employed below. If  $m$  is the mass of a molecule in grammes, then  $\mu = Lm$ ; in particular, for atomic hydrogen ( $\mu$  almost exactly = 1),  $Lm_{\text{H}} = 1$ . If further  $n$  is the number of molecules in the unit of volume,  $N$  that in the volume  $V$ , and if  $\nu$  is the number of moles in the volume  $V$ , then we have  $\nu L = nV = N$ . Finally, we denote by  $\rho = nm$  the density of the gas, and by  $v_s = 1/\rho$  its specific volume.

## 2. Fundamental Assumptions of the Kinetic Theory of Gases.

After these prefatory remarks on the atomic theory in chemistry, we now proceed to the *kinetic theory of gases* (Herapath, 1821; Krönig, 1856). Considering the enormous number of particles in unit volume of a gas, it would of course be a perfectly hopeless undertaking to attempt to describe the state of the gas by specifying the position and velocity of every individual particle. As in all phenomena of matter in bulk, we must here have recourse to *statistics*. But the statistics now to be used is of a somewhat different kind from that which we are acquainted with in ordinary life. There, the statistical method consists in recording a large number of events which have occurred, and in drawing conclusions from the numerical data so obtained. Thus, mortality statistics answers the question of how much more probable it is that a man will die at 60 than at 20 years of age; for this purpose we count the number of

cases of deaths of men at these ages over a long period, and take the respective numbers found for the two ages as proportional to the probabilities required.

If we propose to apply the statistical method to the theory of gases, the method of procedure must be essentially different; for an enumeration of the molecules which, for example, occupy a given element of volume at a given moment is utterly impossible. We must therefore proceed indirectly, first introducing assumptions which appear plausible, and then building up the theory on these as foundation. As with every scientific theory, the final warrant for the correctness of the assumptions is the agreement of their logical consequences with experience.

We may wish, for example, to know the probability of finding a gas molecule at a definite spot in the box within which we suppose the gas to have been enclosed. If no external forces act on the molecules, we shall be unable to give any reason why a particle of gas should be at one place in the box rather than at another. Similarly, in this case there is no assignable reason why a particle of the gas should move in one direction rather than in another. We therefore introduce the following hypothesis, the *principle of molecular chaos*: For the molecules of gas in a closed box, in the absence of external forces, *all positions in the box and all directions of velocity are equally probable.*

In the kinetic theory of gases we shall only have to do with mean values, such as time averages, space averages, mean values over all directions, and so on. Individual values entirely elude observation. If  $n_a$  denotes the number of molecules per unit volume with a definite property  $a$ , e.g. with a velocity of definite magnitude, or with a definite  $x$ -component of velocity, then by the mean value of  $a$  we understand the quantity  $\bar{a}$ , where

$$\bar{a} = \frac{\sum n_a a}{\sum n_a}, \quad \text{or} \quad n\bar{a} = \sum n_a a,$$

$n$  standing for  $\sum n_a$ , the number of molecules per cubic centimetre. If, for example, we suppose the velocity of each molecule to be represented by a vector  $v$  with the components  $\xi$ ,  $\eta$ ,  $\zeta$ , and therefore the magnitude  $v = \sqrt{(\xi^2 + \eta^2 + \zeta^2)}$ , and we wish to find the mean value  $\bar{\xi}$  (for molecules with the velocity  $v$ ), then by the principle of molecular chaos with respect to directions of motion, exactly as many molecules of the gas will have a velocity component  $+\xi$  as a component  $-\xi$ ; the mean value  $\bar{\xi}$  must therefore vanish. A value of  $\bar{\xi}$

differing from zero would imply a mean motion of the whole gas in the direction concerned, with this mean velocity.

On the contrary,  $\bar{\xi}^2$  is not equal to zero. From symmetry, we have

$$\bar{\xi}^2 = \bar{\eta}^2 = \bar{\zeta}^2.$$

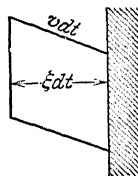
If we take the mean value over all directions, keeping  $v$  fixed, then it follows from  $v^2 = \xi^2 + \eta^2 + \zeta^2$  by taking the mean values that

$$v^2 = \bar{\xi}^2 + \bar{\eta}^2 + \bar{\zeta}^2 = 3\bar{\xi}^2, \quad \text{or} \quad \bar{\xi}^2 = \bar{\eta}^2 = \bar{\zeta}^2 = \frac{v^2}{3}.$$

### 3. Calculation of the Pressure of a Gas.

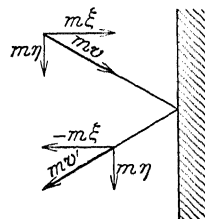
With these ideas before us, we are already in a position to calculate the *gas pressure*  $p$  as the force acting on unit area (D Bernoulli, 1738; Krönig, 1856; Clausius, 1857). According to the kinetic theory of gases, this force is equal to the change of momentum of the molecules striking unit of area of the wall per second. Take the  $x$ -axis at right angles

Fig. 1.—Number of collisions between the wall and molecules having the velocity  $v$ , in the element of time  $dt$ ; it is equal to the number of molecules at a definite moment in the oblique cylinder of height  $\xi dt$  and base the element of area struck.



to the wall. If  $n_v$  denotes the number of molecules in a cubic centimetre which possess the velocity  $v$ , then in the small time  $dt$  a given square centimetre of the wall is struck by all those molecules which at the beginning of the time element  $dt$  were within the oblique cylinder with edge  $v dt$  standing on the square centimetre of the wall (fig. 1). Since the height of this cylinder is  $\xi dt$ ,

Fig. 2.—Momentum diagram for elastic collision of a molecule with the wall; the energy, the magnitude of the momentum and the component momentum parallel to the wall are not changed by the collision, but the component perpendicular to the wall has its sign reversed; momentum  $2m\xi$  is therefore communicated to the wall.



its volume is also  $\xi dt$ ; the number of molecules in it is therefore  $n_v \xi dt$ . The area of the wall considered is struck per unit time by  $n_v \xi$  molecules with the velocity  $v$ .

If we think of the molecules as billiard balls, every molecule when it strikes the wall has its momentum changed by  $2m\xi$  perpendicular to the wall; the component of momentum parallel to the wall is not altered (fig. 2). The molecules considered therefore contribute to the

total pressure  $p$  the part  $2m\xi^2 n_v$ . We sum these values first over all directions of incidence, that is to say, over a hemisphere, keeping  $v$  fixed.

The sum in the present case is equal to half the sum over a complete sphere, so that we have

$$2m\Sigma\xi^2 n_v = 2m \cdot \frac{1}{2} \overline{\xi^2} n_v = \frac{m}{3} v^2 n_v;$$

where  $n_v$  is the number of molecules in a cubic centimetre with a velocity of magnitude  $v$ . If we now sum further over all *magnitudes* of the velocity, we find for the total pressure

$$p = \frac{m}{3} \Sigma n_v v^2 = \frac{m}{3} n \overline{v^2}.$$

If  $V$  is the total volume of the gas, and  $N$  the total number of molecules, it follows from this equation on multiplication by  $V$ , since  $nV = \nu L = N$ , that

$$Vp = N \frac{m}{3} \overline{v^2} = \frac{2}{3} \nu U,$$

where we have put

$$U = L \frac{m}{2} \overline{v^2}.$$

Clearly  $U$  denotes the mean kinetic energy per mole, and in monatomic gases is identical with the total energy of the molecules in a mole. In polyatomic molecules the relations are more complicated, on account of the occurrence of rotations of the molecules, and vibrations of the atoms within a molecule. It can be shown, however, that the preceding formula for the gas pressure holds in this case also;  $U$  denotes as before the mean kinetic energy of the translational motion of the molecules, per mole, but it is no longer the same as the total energy.

#### 4. Temperature of a Gas.

From the kinetic theory of gases, without a knowledge of the law of distribution of velocities (i.e. of the way in which the number  $n_v$  depends on  $v$ ), we have found that the product of pressure and volume is a function only of the mean kinetic energy of the gas. But we have also an empirical law, the *law of Boyle* (1660) and *Mariotte* (1676), viz.: at constant temperature the product of the pressure and volume of an ideal gas is constant. We must conclude from this that  $U$ , the mean kinetic energy per mole, depends only on the temperature of the gas.



In the kinetic theory of gases, the concept of temperature is primarily a foreign element, since in fact the individual molecules are characterized by their velocity alone. But it suggests itself that we should define the absolute gas-temperature  $T$  in terms of the mean kinetic energy. This is usually done in accordance with the equation

$$\frac{m}{2} \overline{\xi^2} = \frac{k}{2} T,$$

where on the left we have the mean kinetic energy of a component of the motion of the centre of inertia of a molecule;  $k$  is called *Boltzmann's constant*. For the total motion of the centre of inertia we have therefore also

$$\frac{m}{2} \overline{v^2} = \frac{3}{2} kT,$$

and, referred to a mole,

$$U = L \frac{m}{2} \overline{v^2} = \frac{3}{2} RT,$$

where we put

$$Lk = R.$$

The justification for this definition of temperature lies in the fact that when we introduce the last expression in the formula for the gas pressure deduced above, we obtain formally the relation which combines the *laws of Boyle-Mariotte and of Gay-Lussac and Charles*:

$$pV = \nu RT.$$

$R$  is called the *absolute gas constant*, and can easily be calculated from measurements of three corresponding values of  $p$ ,  $V$ , and  $T$ . Its value is

$$R = 8.313 \times 10^7 \text{ erg degree}^{-1} \text{ mole}^{-1} = 1.986 \text{ cal degree}^{-1} \text{ mole}^{-1}.$$

We refrain from entering here upon a thoroughgoing discussion of the preceding definition of temperature from the thermodynamic and axiomatic point of view (a complete treatment for the generalized statistics introduced by the quantum theory is given in Appendix XXIX, p. 336), and merely add a brief remark on the units in which temperature is measured.

We use the phrase *ideal gas* if the product  $pV$  is constant at constant temperature; for low pressures this is true for every gas. Deviations from the ideal character of the gas occur when the density of the gas becomes so great that the mean distance between two gas mole-

cules is comparable with the molecular diameter. If we employ such an ideal gas as thermometric substance, the centigrade scale is defined as follows. Let  $(pV)_f$  be the value of  $pV$  for the gas, when it is brought into contact with melting ice, and  $(pV)_b$  its value for contact with boiling water;\* then the temperature of the gas, for the general case of any value of  $pV$ , is defined according to the centigrade scale by

$$t = 100 \frac{pV - (pV)_f}{(pV)_b - (pV)_f}.$$

We see at once that, with this definition, the temperature of melting ice is  $0^\circ \text{C}$ . and that of boiling water  $100^\circ \text{C}$ .

The change from the centigrade scale to the absolute temperature (Lord Kelvin, 1854) scale, which we have indicated above by the symbol  $T$ , is made as follows. It has been established experimentally that at constant pressure an ideal gas expands by  $1/273$  of its volume at  $0^\circ \text{C}$ . when its temperature is raised by  $1^\circ \text{C}$ ., so that we have, e.g.,

$$\frac{(pV)_b}{(pV)_f} = \frac{T_b}{T_f} = 1 + \frac{100}{273} = \frac{373}{273}.$$

If we retain the unit of the centigrade scale in the absolute scale also,  $T_b$  and  $T_f$  must differ by  $100^\circ$ . It follows that melting ice ( $t = 0^\circ \text{C}$ .) has the absolute temperature  $T_f = 273^\circ$ , and boiling water ( $t = 100^\circ \text{C}$ .) the absolute temperature  $T_b = 373^\circ$ . The zero of the absolute temperature scale lies therefore at  $-273^\circ \text{C}$ .

It may be remarked that the absolute temperature scale is sometimes called the *Kelvin scale*, and to distinguish it from the ordinary centigrade scale is denoted by  $K$ .

## 5. Specific Heat.

The *specific heat* (referred to 1 mole) of a substance is given by the energy which must be supplied to the substance to raise its temperature by  $1^\circ$ . For a monatomic gas, it follows immediately from this definition that the specific heat at constant volume is

$$c_v = \frac{dU}{dT} = \frac{3}{2}R.$$

If heat energy is supplied, but with the pressure kept constant instead of the volume, the gas expands and so does work against the external

\*  $f$  = freezing-point,  $b$  = boiling-point.

pressure (which of course in the case of equilibrium must be equal and opposite to the pressure of the gas), the work being

$$p \Delta V = R \Delta T, \quad \text{i.e. } = R \text{ for } \Delta T = 1^\circ.$$

Thus  $R$  is that portion of the specific heat which corresponds to the work of expansion; if we add to this work the specific heat at constant volume  $\frac{3}{2}R$ , we obtain for the total specific heat at constant pressure

$$c_p = \frac{5}{2}R.$$

The ratio  $c_p/c_v$  is generally denoted by  $\gamma$ . Hence, for a monatomic gas we have the relation (Clausius, 1857)

$$\gamma = \frac{c_p}{c_v} = \frac{5}{3} = 1.667,$$

which agrees well with observation, especially for the inert gases.

In polyatomic molecules, besides the three degrees of freedom of the translational motion, other degrees of freedom come in, which correspond to the rotations and vibrations of the molecules and which, when energy is supplied to the gas, can also take up a part of the energy. Now there is a general theorem, the *theorem of equipartition* (Clausius, 1857; Maxwell, 1860), according to which the specific heat is the product of the number of kinetic degrees of freedom by  $k/2$  referred to a molecule, or by  $R/2$  referred to 1 mole. For example, a diatomic molecule, considered as rigid (dumb-bell model), possesses essentially two rotational degrees of freedom. In counting the number of degrees of freedom, the degree of freedom corresponding to rotation about the axis joining the two atoms is to be ignored. For idealized atoms concentrated at a point, this is obviously correct; when we take account of the extension of the atoms in space, we encounter a conceptual difficulty, which it requires the quantum theory to clear up (see Chap. VII, § 2, p. 201). In this case we therefore have  $c_v = \frac{5}{2}R$ ,  $c_p = \frac{7}{2}R$ , and consequently  $\gamma = c_p/c_v = \frac{7}{5} = 1.4$ . These values have in fact been observed, e.g. in molecular oxygen.

## 6. Law of Distribution of Energy and Velocity.

As the next step in the development of the kinetic theory of gases, we proceed to consider the law of distribution of energy or velocity in a gas, i.e., in particular, the law of dependence of the quantity  $n_v$ , employed above, on the velocity. While up to this point a few simple ideas have been sufficient for our purpose, we must now definitely call to our aid the statistical methods of the Calculus of Probabilities.

To begin with a simple example, let us in the first place, without troubling about velocity, consider the question of the number of molecules which on the average are to be found in a definite volume element  $\omega$ . An exact determination of this number for a given moment, apart from the fact that it is inherently impossible, would be of little use, since the number changes every moment on account of the motion of the molecules; we are therefore concerned only with the average number of molecules. This we determine as follows. Imagine the whole box containing the gas to be divided up into separate cells of sizes  $\omega_1, \omega_2, \dots, \omega_z$  (as an aid to the imagination we may picture everything as two-dimensional, see fig. 3), and "throw" the molecules, which we may think of as little balls, at random into this system of cells. We then note that a definite number  $n_1$  of balls has fallen into the first cell, another number  $n_2$  into the second cell, and so on. At a second trial, we shall perhaps find other values for the number of balls in the various cells. If we repeat the experiment a great many times, we shall find that a given distribution, defined by the numbers

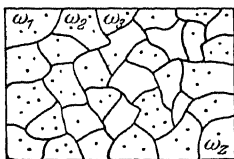


Fig. 3.—Division into cells for the purpose of finding the most probable density of distribution of the molecules of a gas.

$n_1, n_2, \dots, n_z$ , occurs not merely once, but many times. For a model consisting of a real box divided up into cells, and real balls, the frequency with which a given distribution appears could be determined by a long series of trials. For a gas and its molecules this is not possible; we have to use instead an arithmetical argument, consisting of a numerical part (concerning the numbers  $n_1, n_2, \dots, n_z$ ) and a geometrical part (concerning the sizes of the cells  $\omega_1, \omega_2, \dots, \omega_z$ ). In this way we shall find the mathematical probability of a given distribution. And, when we take all possible distributions into account, there is one, for which this probability is greatest, the most probable distribution. Because of the very great number of the molecules in the cubic centimetre, this maximum is overwhelmingly sharp, so that the probability of any distribution deviating essentially from it is negligibly small. It is therefore to be expected that the most probable distribution represents the average state.

At this point, however, a question of principle must be considered. The following difficulty presents itself. If we knew the exact position and velocity of every particle of the gas at a definite moment, the further course of the motion would be completely defined, for the behaviour of the gas is then rigorously determined by the laws of

mechanics, and appears from the very outset to have nothing to do with the laws of probability. If we assume that at time  $t = 0$  the positions and velocities of the molecules are distributed according to any statistical law, we are not entitled without more ado to expect that at any later time  $t$  the state of the gas will be determined by the play of probabilities, independently of the initial state assumed. It would be quite conceivable that, for the initial state chosen, all the molecules would as a consequence of the laws of mechanics be found at the moment  $t$  at a definite corner of the box. In order that it may be possible to apply statistics at all, it must be stipulated that there is no coupling of the states at different instants. We must suppose that the collisions, conditioned by the laws of mechanics, by their enormous number completely efface the "memory" of the initial state after only a (macroscopically) short interval of time. Further, it has to be borne in mind that measurements take time; what we determine is not the micro-state at the instant  $t$  at all, but its mean value over a considerable period. It is assumed that the time-mean values so determined are independent of the period chosen, and that they agree with the mean values ascertained by considerations of probability alone from the most probable state defined above.

Although this hypothesis, the so-called quasi-ergodic hypothesis, is very plausible, its rigorous proof presents difficulties which so far have proved insurmountable. Recently (1932), however, the mathematicians von Neumann and Birkhoff have proved a theorem which is virtually equivalent to the quasi-ergodic theorem. According to the latest ideas in theoretical physics, it is true, the problem of a rigorous proof of the ergodic hypothesis has become unimportant, inasmuch as it is now meaningless, as will be explained later, to give information about the exact position of the individual molecules. To repeat: the hypothesis asserts that even when the initial state is arbitrary a stationary state is in time attained, owing to the collisions of the molecules with one another and with the wall, and that *this state is the same as the state of greatest probability as defined above*. It is assumed that the walls of the box are "rough", so that they do not act as perfect reflectors.

We return now to the calculation of the most probable distribution of the molecules in the individual cells of the box. A definite distribution is described by the numbers in the various cells  $n_1, n_2, \dots, n_z$ ; their sum is of course equal to the number of particles of gas in the box:

$$n_1 + n_2 + \dots + n_z = n.$$

We denote the ratio of the size of the cell  $\omega_k$  to the whole volume  $\omega$  of the box by  $g_k = \omega_k/\omega$ ; we have then

$$g_1 + g_2 + \dots + g_z = 1.$$

How often will this definite distribution actually occur? It is clear in the first place that we obtain the same distribution, if we permute the individual molecules among themselves; the number of these permutations is  $n!$ . Here, however, we are including the cases in which the molecules in a given cell are permuted with one another; since these permutations do not represent new possibilities of realization of the distribution considered, we must divide  $n!$  by the number  $n_1!$  of the permutations within the first cell, and so on, and thus obtain for the number of possibilities of realization

$$\frac{n!}{n_1! n_2! \dots n_z!}.$$

To get the probability of this distribution, we have still to multiply this number by the a priori probability of the distribution, which is  $g_1^{n_1} g_2^{n_2} \dots g_z^{n_z}$ , since the a priori probability that a particle should fall into the first cell is  $g_1$ , and therefore the same probability for  $n_1$  particles is  $g_1^{n_1}$ , and so on. The probability of our distribution given by the numbers  $n_1, n_2, \dots$  therefore becomes

$$W = \frac{n!}{n_1! n_2! \dots n_z!} g_1^{n_1} g_2^{n_2} \dots g_z^{n_z}.$$

To verify that in our calculation we have actually included all the possibilities of realization, we can find the sum of the probabilities for all possible distributions, which of course must be 1, since it is a certainty that one of the distributions is realized. We therefore form the sum over all distributions  $n_1, n_2, \dots$ , for which  $n_1 + n_2 + \dots = n$ . This sum is easily found by the polynomial theorem, which gives

$$\sum_{n_1, n_2, \dots} \frac{n!}{n_1! n_2! \dots n_z!} g_1^{n_1} g_2^{n_2} \dots g_z^{n_z} = (g_1 + g_2 + \dots + g_z)^n = 1.$$

We now transform the above formula for the probability of a definite distribution by means of Stirling's theorem, which for large values of  $n$  gives

$$\log n! = n(\log n - 1).$$

We then obtain from the expression for  $W$ , by taking logarithms,

$$\log W = \text{const.} + n_1 \log \frac{g_1}{n_1} + n_2 \log \frac{g_2}{n_2} + \dots$$

To find the most probable distribution, we must calculate the maximum of  $\log W$  for a variation of the numbers  $n_1, n_2, \dots$ , subject to the subsidiary condition  $n_1 + n_2 + \dots + n_z = n$ . Since  $n_1, n_2, \dots$  are very large we can treat them as continuous variables. By the method of Lagrange's multipliers we obtain as the equations defining this maximum

$$\frac{\partial \log W}{\partial n_1} = \log \frac{g_1}{n_1} - 1 = \lambda, \quad \frac{\partial \log W}{\partial n_2} = \log \frac{g_2}{n_2} - 1 = \lambda, \dots,$$

where  $\lambda$  is a constant, whose value is determined by the equation  $n_1 + n_2 + \dots = 1$ . It follows that

$$\frac{g_1}{n_1} = \frac{g_2}{n_2} = \dots = e^{\lambda+1} = \text{const.};$$

$$n_1 = ng_1 = n \frac{\omega_1}{\omega}, \quad n_2 = ng_2 = n \frac{\omega_2}{\omega}, \dots$$

Hence the numbers  $n_1, n_2, \dots$  in the individual cells are proportional to the size of the cells; we therefore have a *uniform distribution of the molecules* over the whole box; the size of the cells does not matter at all.

While the result just deduced for the distribution of the density of the molecules was to be expected from the start, the same method, applied to the distribution of the velocity of the molecules, leads to a new result. The calculations in this case are exactly analogous to those above. We construct a "velocity space" by drawing lines from a fixed point as origin, representing as vectors the velocities of the individual molecules in magnitude and direction. We then investigate the distribution of the ends of these vectors in the velocity space. In this case as before we can make a partition into cells, and consider the question of the number of vectors whose ends fall in a definite cell. There is, however, one essential difference as compared with the former case, in that there are now *two subsidiary conditions*, viz. besides the condition

$$n_1 + n_2 + \dots + n_z = n$$

for the total number of particles, an additional condition for the total energy  $E$  of the gas, viz.

$$n_1 \epsilon_1 + n_2 \epsilon_2 + \dots + n_z \epsilon_z = E,$$

where  $\epsilon_l$  denotes the energy of a molecule, whose velocity vector points into the cell  $l$ . Taking account of these two subsidiary conditions, we

find for the maximum of the probability ( $\lambda$  and  $\beta$  being the two Lagrangian multipliers)

$$\frac{\partial \log W}{\partial n_l} = \log \frac{g_l}{n_l} - 1 = \lambda + \beta \epsilon_l \quad (l=1, 2, \dots, z).$$

This leads to *Boltzmann's law of distribution* (1896)

$$n_l = g_l e^{-1-\lambda-\beta\epsilon_l} = g_l A e^{-\beta\epsilon_l},$$

where  $A$  and  $\beta$  are two constants, which have to be found from the two subsidiary conditions. Thus the expression for the number  $n_l$  corresponding to the cell  $l$  essentially involves the energy belonging to this cell as well as the size  $g_l$  of the cell, and that in such a way that

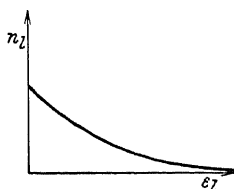


Fig. 4.—Boltzmann's law of distribution: if the cells are of equal size, those with greater energy are more sparsely filled than those with smaller energy.

among cells of equal size one with greater energy is not so well filled as one with smaller energy; the fall in the value of the number  $n_l$  with increasing energy obeys an exponential law (fig. 4).

We shall now apply *Boltzmann's law of distribution* to the *special case of a monatomic gas*. Here the energy is given by

$$\epsilon = \frac{m}{2} v^2 = \frac{m}{2} (\xi^2 + \eta^2 + \zeta^2).$$

Position in the velocity space is uniquely defined by the three components  $\xi, \eta, \zeta$ . From their meaning the cells  $g$  are finite. From the macroscopic standpoint we can consider them as infinitesimal, and denote them by  $d\xi d\eta d\zeta$ . In the formation of mean values we can then replace the sums by integrals:

$$\sum_l g_l \dots \rightarrow \iiint d\xi d\eta d\zeta \dots,$$

the limits of integration being  $-\infty$  and  $+\infty$  in each case. Since further we are interested in the mean values of  $v, v^2, \dots$  only ( $\bar{\xi}, \bar{\eta}, \bar{\zeta}$  vanish from symmetry, and  $\bar{\xi}^2, \bar{\eta}^2, \bar{\zeta}^2 = \frac{1}{3}\bar{v}^2$ ), the integrand in each of the integrals to be calculated depends on  $v$  only, which suggests that we should introduce polar co-ordinates in the velocity space, with  $v$  as radius. Integration with respect to the polar angles can be effected



at once, giving the factor  $4\pi$ , the area of the unit sphere. We have therefore

$$\int \int \int d\xi d\eta d\zeta \dots = 4\pi \int_0^\infty v^2 dv \dots$$

The total number  $n$  of the molecules is obtained from Boltzmann's law by calculating the integral

$$n = 4\pi A \int_0^\infty v^2 e^{-\beta \cdot \frac{1}{2}mv^2} dv,$$

while for the total energy we have

$$E = 4\pi A \int_0^\infty \frac{m}{2} v^4 e^{-\beta \cdot \frac{1}{2}mv^2} dv.$$

These two equations determine uniquely the two constants  $A$  and  $\beta$ , which up to this point were unknown. The integrals are readily evaluated (see Appendix I, p. 259); we find the relations

$$n = A \sqrt{\left(\frac{\pi}{\lambda}\right)^3}, \quad \left(\lambda = \frac{\beta m}{2}\right)$$

$$E = \frac{3}{4} m A \sqrt{\left(\frac{\pi^3}{\lambda^5}\right)} = \frac{3}{4} \frac{mn}{\lambda} = \frac{3}{2} \frac{n}{\beta}.$$

But we have seen (p. 7) that on the average a molecule, corresponding to its three translational freedoms, possesses the kinetic energy  $\frac{3}{2}kT$ . Since the whole kinetic energy of the gas is therefore  $\frac{3}{2}nkJT$ , it follows that

$$\beta = \frac{1}{kT}.$$

The constants in Boltzmann's law are accordingly expressed in terms of the number of molecules in the gas and their absolute temperature.

We inquire next how many molecules have a velocity between  $v$  and  $v + dv$ . This number  $n_v dv$  is clearly given by the integrand of the above integral for  $n$ :

$$n_v = 4\pi A e^{-\lambda v^2} v^2 = 4\pi n \left(\frac{m}{2\pi kT}\right)^{3/2} e^{-mv^2/2kT} v^2.$$

This relation is known as *Maxwell's law of distribution of velocity* (1860); the graph of  $n_v$  as a function of  $v$  is shown in fig. 5.

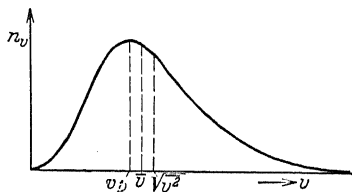


Fig. 5.—Maxwell's law of distribution of velocities, showing the most probable velocity ( $v_p$ ), the mean velocity ( $v$ ), and the root-mean-square velocity  $\sqrt{v^2}$ .

To obtain some idea of the order of magnitude of the velocities of gas molecules, we can calculate from the law of distribution the most probable velocity  $v_p$ , or the average velocity  $\bar{v}$ , or other such mean values (see Appendix I, p. 259); we find, e.g., for the most probable velocity the value  $v_p = \sqrt{(2RT/\mu)}$ , so that, e.g., for molecular hydrogen ( $\mu = 2$ ) at  $0^\circ \text{C}$ . ( $T = 273^\circ \text{K}$ .)

$$v_p = 15.06 \times 10^4 \text{ cm./sec.}$$

An experimental test of Maxwell's law of distribution may be carried out as follows. Let a furnace O contain a gas at a definite temperature  $T$ . In the wall of the furnace let there be an opening through which



the gas molecules can escape into a space which is highly evacuated (fig. 6). The escaping molecules fly on in a straight line with the velocity which they possessed in the furnace at the moment of their exit through the opening. By means of a diaphragm system behind the opening, the stream of molecules issuing in all directions may be cut down to a *molecular ray* (Dunoyer, 1911). The distribution of velocity in the ray, or beam, can be measured directly by various methods, the

Fig. 6.—Diagrammatic representation of the production of a molecular ray. The furnace O, which contains the gas, is heated from outside.

most important of which will be described immediately. When deducing from this distribution the distribution of velocity in the enclosed gas, we must, however, take the fact into account that rapid molecules are relatively more numerous in the beam than in the gas. For the beam consists of all molecules which issue through the opening per unit time; the number of these is (§ 3, p. 6) proportional to  $n_v v dv$ , whilst  $n_v dv$  represents the number of particles having the same range of velocity in the gas. The two distributions therefore differ from each other by the factor  $v$ .

A direct method (Stern, 1920) of measuring the velocity in the beam depends on the following principle. A beam, consisting, e.g., of silver atoms, can easily be demonstrated by setting up in the path of the beam a glass plate, on which the silver is deposited. If the apparatus is then rotated about an axis perpendicular to the path of the beam, the molecules are no longer deposited at the same spot on the glass plate as when the tube was at rest, but at a greater or smaller distance from this point according to their velocity, since, during the time of flight of the molecules from the furnace to the glass plate, the tube and with it the glass plate has turned. The distribution of velocity

in the beam is therefore found immediately from the plate, by measuring the intensity of the deposit at various distances from the original point.

A more recent method makes use of the same principle as was used by Fizeau (1849) to measure the velocity of light, viz. that of coupled rotating toothed wheels. Further details of the method need not be gone into here. The measurements, as carried out mainly by Stern and his pupils, showed that the distribution of velocity among the molecules in the furnace does actually satisfy Maxwell's law.

Another method depends on the Doppler effect (1842). If a molecule, which in the state of rest emits a definite frequency  $\nu_0$ , moves towards the observer with the relative velocity component  $v_r$ , the light appears to him to be displaced in the direction of higher frequencies by the factor  $1 + v_r/c$ ; while for the opposite relative motion the corresponding factor is  $1 - v_r/c$ . When the light which is emitted by a luminous gas is spectrally decomposed, then besides the frequency  $\nu_0$  there occur also all frequencies which are derived from  $\nu_0$  in consequence of the Doppler effect due to the motion of the molecules, and that with an intensity which is given by the number of molecules with a definite component of velocity relative to the observer. Each spectral line has therefore a finite width, and the intensity distribution over its range gives a picture of Maxwell's law.

## 7. Free Path.

We have spoken above of a molecular ray. This consists of molecules which have passed through the diaphragm system and fly on as a beam through the evacuated space. A necessary condition is a *high vacuum*. If, however, there is still a residue of gas in the space (with molecules either of the same or of a different sort), then some of the molecules in the beam, more or fewer according to the pressure, will collide with the gas molecules and be scattered in the process; consequently the molecular beam will be weakened, clearly, in fact, exponentially with the distance travelled, since the number of scattering processes is proportional to the number of molecules present in the beam. If then we denote by  $n(s)$  the number of beam molecules which, after traversing a distance  $s$  from the opening in the furnace, pass through a plane at right angles to the beam per unit time, we have a law of the form

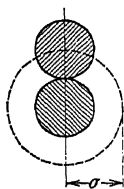
$$n(s) = n(0)e^{-s/l}.$$

Here  $l$  is a quantity (to be determined later) of the dimensions of a length; a little consideration shows that this quantity is equal to the distance which on the average a molecule of the beam covers before

it comes into collision with a molecule of the gas outside the beam. For, putting  $1/l = \beta$ , we have for this average

$$\begin{aligned}\bar{n} &= \frac{\int_0^\infty sn(s) ds}{\int_0^\infty n(s) ds} = \frac{\int_0^\infty se^{-\beta s} ds}{\int_0^\infty e^{-\beta s} ds} \\ &= -\frac{d}{d\beta} \left( \log \int_0^\infty e^{-\beta s} ds \right) = -\frac{d}{d\beta} \left( \log \frac{1}{\beta} \right) \\ &= \frac{d}{d\beta} (\log \beta) = \frac{1}{\beta} = l.\end{aligned}$$

We therefore call  $l$  "the mean free path" of the ray in the gas (Clausius, 1858). The present author, together with E. Bormann, has shown (1921) how it can be determined on the basis of its definition by the exponential law, by measuring the attenuation of a beam of silver atoms on passage through the gas which is at rest (air). The more important case is that in which beam and gas consist of molecules of the same kind; the mean free path  $l$  is then a property of this gas.



The question of the quantities on which the mean free

Fig. 7.—Gas-kinetic effective cross-section; in a collision the centres of mass of the two equal molecules cannot come nearer than the distance  $\sigma$  ( $\sigma$  = diameter of a molecule).

path must depend may be examined theoretically. Clearly it is a matter of the number of collisions with the molecules of the gas, in which a definite molecule takes part during its passage through the gas. We may regard the other molecules of the gas as at rest; to take their motion into account does not make any essential difference. We shall regard the molecules as spheres of diameter  $\sigma$ , and have to answer the question of how many collisions such a moving sphere is involved in, during its flight through the gas, which consists of spheres at rest. Since a collision occurs whenever the centre of the moving molecule, in the course of its flight, comes nearer the centre of a molecule at rest than the distance  $\sigma$ , we may also obtain the number of these collisions by considering spheres of radius  $\sigma$  at rest, and a moving point (fig. 7). We have therefore the same problem before us as when a man fires a gun in a wood, and considers the question of the number of trees he has struck. The number is clearly proportional to the thickness of an individual tree, and to the number per unit area; its reciprocal determines the mean range of the bullets. In the case of the gas also, the collision

number must be proportional to the number  $n$  of the molecules of gas per unit volume and to their (gas-kinetic) cross-section  $\pi\sigma^2$ . Since the mean free path is inversely proportional to this collision number, it is proportional to

$$\frac{1}{n\pi\sigma^2} = \frac{V}{\pi L\sigma^2},$$

where  $V$  is the volume per mole.

Measurement of the mean free path  $l$  therefore throws light upon the value of the product  $L\sigma^2$ . A direct method of determining  $l$  (for a beam of molecules of a foreign element) has already been given above. As to indirect methods, we must mention first a method, given by Maxwell (1860), depending on the conduction of heat in the gas (Appendix II, p. 261). If the gas molecules did not collide, a rise in temperature anywhere in the gas, that is to say, an increase in the kinetic energy of the particles, would be propagated through the gas with the great velocity of the molecules, say a thousand metres per second; experimentally, however, gases are found to be relatively poor conductors of heat. The reason for this is that a gas molecule can fly for only a relatively short distance, of the order of magnitude of the free path, before it collides with another particle of gas, and thus not only changes its direction of motion, but also gives up a part of its kinetic energy to the particle it has struck. Other methods for the determination of  $l$  depend on *viscosity* and on *diffusion* (Appendix II, p. 263). The latter method lends itself well to visual demonstration; if, for example, we let chlorine gas diffuse into air, then, since the colour of the chlorine makes it visible, we can observe directly how slowly the diffusion goes on.

From these experiments, taken together, we find for the order of magnitude of the mean free path at a pressure of 1 atmosphere,  $l \sim 10^{-6}$  cm.; at a pressure of  $10^{-4}$  mm. of mercury, which corresponds to the normal X-ray vacuum,  $l \sim 10$  cm.

## 8. Determination of Avogadro's Number.

As has already been remarked above, when the free path is known so is the product  $L\sigma^2$ , i.e. the product of the square of the molecular diameter and Avogadro's (Loschmidt's) number. To determine  $\sigma$  and  $L$  separately, we need a second relation between them. Such a relation is given, at least in respect of order of magnitude, by the *molar volume* of the solid body. In the solid state of aggregation, it is reasonable to suppose that the molecules are in the state of densest

packing, so that the volume taken up by 1 mole—up to a factor of the order of magnitude 1—is given by the product of the number of molecules per mole and the volume taken up by one molecule, i.e. by  $L\sigma^3$ . From  $L\sigma^2$  and  $L\sigma^3$  we can now determine  $L$  and  $\sigma$ , so obtaining

$$L \sim 10^{23} \text{ mol}^{-1}, \sigma \sim 10^{-8} \text{ cm.}$$

Moreover, the influence of the “proper volume” of the molecules is shown not only in the most condensed state (the solid body), but even in the gaseous state by deviations from the law of ideal gases

$$pV = RT.$$

Thus, if the volume  $V$  of a definite mass of gas becomes so reduced that the proper volume of the molecules is actually comparable with  $V$ , then the free volume at the disposal of the individual molecule is smaller than  $V$ , and we obtain the equation of state

$$p(V - b) = RT.$$

The exact calculation gives for  $b$  four times the proper volume of the molecules. (For spherical particles of diameter  $\sigma$ ,

$$b = 4L \cdot \frac{4\pi}{3} (\sigma/2)^3 = \frac{2\pi}{3} L\sigma^3;$$

see Appendix III, p. 265). In dense gases, however, there are also other deviations from the ideal equation of state, which are due to the *cohesion* of the molecules, and have the general effect of making the pressure, for given  $T$  and  $V$ , smaller than it is according to the formula  $pV = RT$ . To represent these circumstances, many equations of state have been proposed; the best known is that of van der Waals (1881),

$$\left(p + \frac{a}{V^2}\right)(V - b) = RT.$$

What chiefly interests us here, is that by determination of the constant  $b$  we again obtain the product  $L\sigma^3$ . At pp. 253, 339 we shall return to the question of the value of  $a$ , which measures the cohesion.

The evaluation of Avogadro's number given above is of course rather rough. We get a more exact method by considering *fluctuation phenomena*. Thus, if we take 1 cm.<sup>3</sup> of a gas, we shall find exactly as many molecules in it as in another cm.<sup>3</sup>, viz., at room temperature, about  $10^{19}$ ; differences of a few hundred molecules in these huge numbers are of course of no moment. It is another matter if we pass to

smaller elements of volume; in a cube of edge  $0.1 \mu$  there are on the average only  $10^4$  molecules, and it is clear that variations of a few hundred molecules are now relatively quite important. If we go on to smaller regions still, we shall finally arrive at volumes which contain only one or two molecules, or none at all. The smaller, therefore, the number of particles considered, the greater will the fluctuations be (Appendix IV, p. 266).

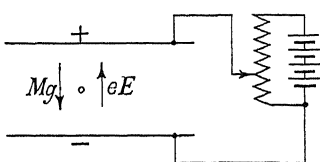
Examples of these fluctuation phenomena are given by the *Brownian movement* (1828), which is observed with microscopic particles (e.g. colloidal solutions, or smoke in air), and manifests itself macroscopically in the oscillations of a mirror suspended by a fine wire; also by the *sedimentation of suspensions*, in which colloidal particles by reason of their weight tend to sink to the bottom of the vessel, but in consequence of collisions with the particles of the solvent are impelled more or less upwards, thus giving rise to a density distribution of the same character as that in the atmosphere, expressed in the barometric height formula. A third example is the *scattering of light in the atmosphere* (Lord Rayleigh, 1871), which causes the colour of the sky. If, in fact, the density of air were the same throughout, then just as in an ideal crystal there would be no scattering of light, since the scattered waves leaving individual molecules would annul each other by interference; the sky would appear black. Scattering is only possible when there are irregularities in the uniformity of distribution, i.e. fluctuations of density; indeed these fluctuations must be so pronounced that they are perceptible within a distance of the order of magnitude of a wave-length. Since the fluctuations in small volumes are greater, short (blue) waves are more strongly scattered than long (red) waves; the sky therefore appears blue.

A much more exact determination of  $L$  depends on its connexion with the *elementary electrical charge*  $e$ , and the electrolytic unit, the *faraday*. In electrolysis, we have the following fundamental law, discovered by Faraday (1833). In the electrolytic separation of 1 mole of any substance, a quantity of electricity is transported equal to 96,540 coulombs; i.e.

$$eL = F = 96,540 \text{ coulombs.}$$

The interpretation is obviously that every ion carries the same elementary charge  $e$  (see § 1, p. 44) and that the total amount of electricity transported is  $e$  times the number of atoms. If then we know  $e$ , we can calculate from it the value of  $L$ , or of  $m_H = 1/L$ . A method for determining  $e$  was devised by Ehrenhaft (1909) and Millikan (1910), and

the latter has developed it to high precision. If electricity actually consists of elementary quanta, the total charge on a body must be an integral multiple of the charge  $e$ . Owing to the smallness of the elementary charge, it is of course difficult to test the truth of this by experiments on macroscopic charged bodies. Such experiments offer little



prospect of success unless the whole charge consists of very few quanta; it actually does so in Millikan's oil-drop

Fig. 8.—Condenser in Millikan's method of determining the elementary charge  $e$ . The weight  $Mg$  is compensated by an electric field  $E$ .

method (fig. 8). The charge on an oil droplet may be determined with sufficient exactness, by bringing it into the field of a condenser whose lines of force are directed vertically upwards. There are then two forces acting on the particle in opposite directions, viz. the electric force upwards, and the force of gravity downwards; the droplet will be in equilibrium if the potential difference applied to the condenser is just so chosen that

$$eE = Mg,$$

where  $e$  and  $M$  are the charge and mass of the oil-drop,  $E$  is the electric field strength, and  $g$  is the acceleration of gravity. The greatest difficulty is the determination of the mass  $M$ . It can be calculated from the density and the radius of the drop, if it is assumed that the density has its normal value. The radius is determined by switching off the electric field, so that the particle falls; and then measuring the velocity of descent, which, in consequence of the viscosity of the surrounding medium, is constant. Inserting this velocity in the formula called Stokes's law, we obtain the radius of the droplet. The experiment, as carried out by Millikan and others, not only proved without ambiguity that the charge on the droplets consists of integral multiples of an elementary quantum, but also allowed an exact determination to be made of this elementary charge  $e$ ; the value found is

$$e = 4.77 \times 10^{-10} \text{ electrostatic units.}$$

By using this result in Faraday's law, we obtain for Avogadro's number the value

$$L = 6.06 \times 10^{23} \text{ molecules per mole.}$$

We may also refer here to various methods for the determination of  $L$ , which make use of the *radiations from radioactive substances*. We may,



e.g., count the number of particles, either by noting the scintillations produced when they strike a screen moistened with zinc sulphide, or with the help of the Geiger counter. If we know the mass of our sample and the decay constant, we can calculate  $L$ .

In conclusion, it may be mentioned in addition that the Boltzmann constant  $k$ , which by definition is the quotient of the gas constant  $R$  by Avogadro's number, can be also measured directly by determining the *spectral distribution of intensity in the radiation emitted by a black body*. The function which expresses the intensity in terms of the frequency and the temperature involves only two universal constants,  $k$  and  $h$ , the first of which is Boltzmann's constant; the second is called Planck's constant, and is the fundamental constant of the quantum theory (Chap. VII, § 1, p. 185).

## CHAPTER II

# Elementary Particles

### 1. Conduction of Electricity in Rarefied Gases.

The development of chemistry and of the kinetic theory of gases has led, as explained in the preceding chapter, to the assumption that matter consists of molecules and atoms. For the *chemist* these particles represent the ultimate constituents of which solids, liquids, and gases are composed, and on which he acts in every reaction produced by purely chemical means.

We now take up this question of the smallest components of matter from the point of view of the *physicist*. The latter has many other phenomena at his command, from the study of which he can collect data upon the structure of matter. Among these phenomena, the processes which occur when electricity passes through rarefied gases have proved to be of special importance.

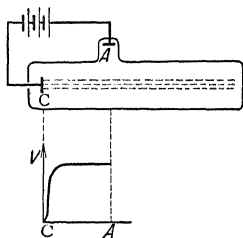
Under normal conditions, a gas is in general a poor conductor of electricity. If, however, the gas is enclosed in a vessel with two electrodes to which a sufficiently high potential is applied, it is found that at pressures of a few millimetres of mercury, a transport of electricity takes place across the gas, showing itself by the flow of an electric current in the external leads to the electrodes; and that at the same time the gas becomes strongly illuminated, a phenomenon which from the theoretical point of view is somewhat complicated, but which in the form of the so-called Geissler tubes is extensively employed in practice, especially for illuminated signs.

If the pressure of the enclosed gas is reduced still further, the illumination disappears (below 0.1 mm. of mercury) almost completely. At very low pressures (about  $10^{-5}$  to  $10^{-6}$  mm.), however, the presence of rays is observed, proceeding from the cathode and producing a fluorescent appearance on the opposite wall of the glass tube. In certain circumstances these rays can be seen directly in the form of

a bluish thread, stretching from the cathode across the tube. The rays are called *cathode rays* (fig. 1).

Their properties can be investigated as follows (Plücker, 1858; Hittorf, J. J. Thomson). If a body is placed in the path of the rays, it is seen to cast a shadow on the fluorescent part of the glass. From the geometrical relations it can be inferred that the rays producing the shadow *spread out in*

Fig. 1.—How the cathode rays arise. The rays proceed from the cathode C, and can be seen as a bluish thread if the gaseous pressure in the tube is not too small. A is the anode. The propagation of the rays in straight lines from the cathode is due to the fact that the potential between C and A, in consequence of the presence of slight residues of gas (space charges) in the tube, does not rise uniformly from C to A, the rise taking place almost entirely in the immediate neighbourhood of the cathode. [The graph of the potential is shown in the lower part of the figure.] The circuit is closed by a weak transport of positive ions.



*straight lines*. It is found also that the occurrence of these rays is associated with a transport of electric charge across the tube. Moreover, the rays can be deflected from their rectilinear path by external electric or magnetic fields, and that in such a way as to permit the inference that the rays consist of rapidly moving, *negatively charged particles*; they are called *electrons*.

The velocity and the specific charge (i.e. the ratio  $e/m$  of charge to mass) of these particles can also be determined. If we set up two wire gratings in the path of the rays and at right angles to their direction (fig. 2), and apply to them a potential difference  $V$ , the electrons in the *longitudinal electric field* between the gratings are accelerated or retarded. The change of speed due to their passage across the field is given by the equation of energy. If  $v_0$  is their velocity *before*,

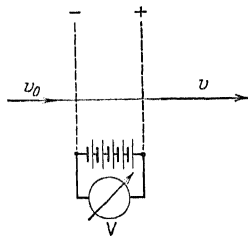


Fig. 2.—Acceleration of electrons by an electric field; the increase of kinetic energy is equal to the potential fall multiplied by the charge  $-e$  of the electron.

and  $v$  their velocity *after* passing through the field, then for the case where the field is in the opposite direction to the velocity so that the negative electrons are accelerated,

$$\frac{m}{2} v_0^2 = \frac{m}{2} v^2 + eV.$$

When the initial velocity is small, we may put  $v_0 \approx 0$ , i.e.

$$\frac{m}{2} v^2 = -eV.$$

If  $V$  is known, we can therefore determine in this way the value of  $(m/e)v^2$ .

The same quantity can also be measured by deflection in a *transverse electric field* (fig. 3). If two condenser plates are set up parallel to the path of the rays and at a distance  $l$  apart, and a potential difference  $V$  is

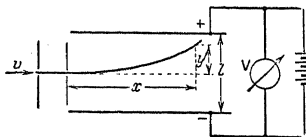


Fig. 3.—Deflection of an electron in a transverse electric field (between the plates of a condenser). The path is a parabola (path of a projectile).

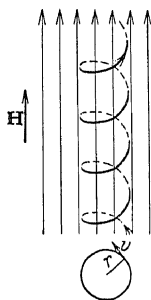
applied to them, a constant deflecting force  $eV/l$  acts on the electrons perpendicular to their original direction of motion. Their path is therefore a parabola, which is defined by the equations

$$x = v_0 t, \quad y = \frac{g}{2} t^2,$$

where  $v_0$  is the initial velocity, and  $g = (e/m)(V/l)$  is the acceleration due to the field. Elimination of the time  $t$  gives

$$y = \frac{e}{2m} \frac{V}{l} \frac{x^2}{v_0^2}, \quad \text{or} \quad \frac{m}{e} v_0^2 = \frac{V}{l} \frac{x^2}{2y}.$$

By measuring the deflection  $y$  of the ray after it has traversed a distance  $x$ , we can therefore again find the value of the quantity  $(m/e)v_0^2$ .



In a *magnetic field*  $H$  we know that a moving charge  $e$  is acted on by a force which is perpendicular both to the direction of the field and to that of the velocity; the magnitude of this so-called *Lorentz force* is given by  $e(v/c)H$ , multiplied by the sine of the angle

Fig. 4.—Motion of an electron in a constant magnetic field  $H$ ; the path is in general a helix, with its axis parallel to the field; in the special case when the electron moves at right angles to  $H$ , the path is a circle.

between the directions of velocity and field ( $c = 300,000$  km./sec. = the velocity of light). The path of the electron is a helix, whose axis is parallel to the field (fig. 4). In the special case when the component velocity parallel to  $H$  vanishes, the helix shrinks into a circle perpendicular to the field. Its radius  $r$  is easily calculated. The Lorentz force  $e(v/c)H$  is directed towards the centre of the circle. It must obviously be equal to the centrifugal force  $mv^2/r$ , so that

$$\frac{e}{c} v H = \frac{mv^2}{r}, \quad \text{or} \quad \frac{m}{e} v = \frac{Hr}{c}.$$

Measurement of the radius  $r$  and of the magnetic field strength therefore enables us to find the value of  $(m/e)v$ .

We thus obtain the *result*, that measurements of deflection give in the electric field  $(m/e)v^2$ , and in the magnetic field  $(m/e)v$ . Hence the values of  $e/m$  and  $v$  can be determined. Actual measurements have shown velocities which, with increasing potential difference across the tube, reach near the velocity of light.

With regard to the measurement of  $e/m$ , exact experiments have shown that this, the specific charge, is not precisely constant, but depends to some extent on the velocity of the electrons (Kaufmann, 1897). This phenomenon is explained by the *theory of relativity* (Appendix V, p. 269). According to Einstein (1905), the value of the charge  $e$  is invariable; the mass, however, is variable, its magnitude in fact depending upon the velocity which it has, relative to the observer who happens to measure it. The electron may have the "rest mass"  $m_0$ , i.e. this is its mass for the case when it is at rest relative to the observer; but if it moves relative to the observer with a velocity  $v$ , it behaves (e.g. in a field of force) as if it possessed the mass

$$m = \frac{m_0}{\sqrt{\left(1 - \frac{v^2}{c^2}\right)}}$$

This assertion of the theory can be tested by deflection experiments on cathode rays; these experiments confirm it completely (see § 1, p. 44).

The result given by the deflection experiments for the limiting value  $e/m_0$ , i.e. for the specific charge of an electron reduced to zero velocity, is

$$\frac{e}{m_0} = 1840 F,$$

where  $F$  denotes Faraday's constant, i.e. represents the quantity of electricity transported in the electrolytic separation of 1 mole. It is given (p. 21) by

$$F = \frac{e}{m_{\text{H}}},$$

where  $m_{\text{H}}$  is the mass of a hydrogen atom. For the rest-mass of the electron we have therefore the relation

$$m_0 = \frac{m_{\text{H}}}{1840} = 9.0 \times 10^{-28} \text{ gm.}$$

## 2. Canal Rays and Anode Rays (Positive Rays).

We have in the cathode rays made the acquaintance of electrically negative charged particles. An obvious question now is: can we not produce positively charged rays also, in the same kind of way as the cathode rays? The answer was given by Goldstein (1886), who succeeded in producing such rays by the following method. If residues of gas are still present in the discharge tube, the electrons on their way from the cathode to the anode will collide with these residual gas molecules and ionize them. The ions thus formed, being positively charged, will be accelerated in the direction towards the cathode, in consequence of the potential difference which exists across the discharge tube. They therefore dash against the cathode, and

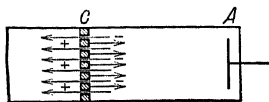


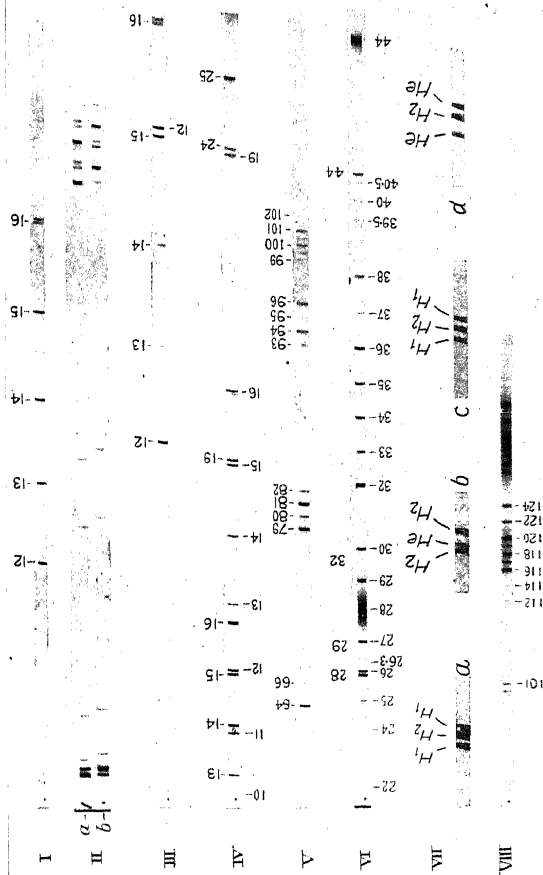
Fig. 5.—Production of canal rays; the positive ions produced in front of the cathode are driven up to the cathode and pass through the canals bored in it.

would of course stick fast in it unless, as Goldstein did, we bored canals through the cathode, which allow the ions free passage (fig. 5). The rays so obtained are called *canal rays*.

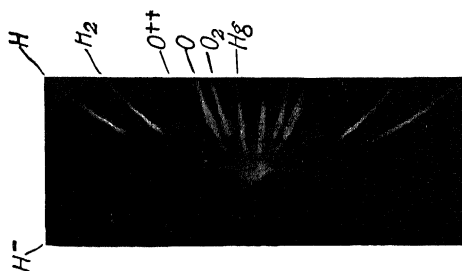
Under certain conditions, positively charged rays issue also from the anode, consisting of ions which have been torn out of the atomic fabric of the anode; rays of this sort are called *anode rays*.

The properties of these *positive rays* can be determined by methods analogous to those used for the cathode rays. From the deflection experiments values are found for the specific charge of these particles, of the order of magnitude of Faraday's number  $F$ . In these rays, therefore, we are concerned with singly or multiply charged atoms or molecules (ions); and in fact the same values are found for the ratios of the masses of these ions, as chemists have found by chemical methods.

For the exact determination of the specific charge, an apparatus was constructed (J. J. Thomson) in which the ions are deflected in an electric field and a magnetic field parallel to it. If a photographic plate is set up perpendicular to the original direction of the rays, the image obtained on the plate is a family of parabolas (fig. 6, Plate I). By a simple calculation (§ 1, p. 26) we can see that the points on a definite parabola arise from a definite set of particles with the same value of  $e/m$ ; and that the individual points of this parabola correspond to different velocities of the particles, in such a way in fact that the marks due to the particles with smaller velocities, and therefore more easily deflected, are farther away from the vertex of the para-



Ch. II, Fig. 8.—Examples of mass-spectra (from Aston's *Isotopes*, by courtesy of the publishers, Edward Arnold & Co.)



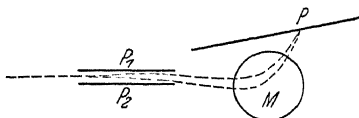
Ch. II, Fig. 6.—Positive ray parabolas (J. J. Thomson). Measurements on these deflection photographs allow the individual parabolas to be assigned to chemical elements and compounds as indicated. The electric and magnetic deflecting fields are parallel to the common axis of the parabolas.





bola. Since a definite parabola corresponds to each value of  $e/m$ , we can easily, from the positions of the individual parabolas, determine by measurement the specific charges of the ions contained in the ray, and accordingly their masses also (since we know their charge, which of course is an integral multiple of the elementary charge).

Fig. 7.—Diagrammatic representation of Aston's mass-spectrograph. The positive ray is deflected, first downwards in an electric field (between the condenser plates  $P_1$  and  $P_2$ ), and then upwards in a magnetic field (indicated by the coil  $M$ ). By suitable dimensioning and arrangement of the apparatus it can be secured (as was shown by Aston and Fowler) that positive rays with the same  $e/m$  but arbitrary velocity are focussed at one and the same point  $P$  of the photographic plate. A mass-spectrum is obtained on the photographic plate; examples are shown in fig. 8, Plate I.



Aston's "mass-spectrograph" (for construction, see fig. 7; for mass-spectra, fig. 8) enables us to determine the mass directly. It has the immediate advantage over chemical methods, that the measurement of mass is made on the individual ion, whereas chemists always measure only the mean value of the mass over a very large number of particles. We shall return to this subject later (p. 37).

### 3. X-rays.

In the year 1895 Röntgen discovered a new kind of ray, distinguished by a penetrating power up to that time unknown. There are still many who remember what a sensation it made when the first photograph of the bones of a living subject was published. The hopes then aroused in the medical profession have been to a large extent fulfilled. But in physics also new paths were opened up by this discovery. The ray-physics characteristic of the present day had begun.

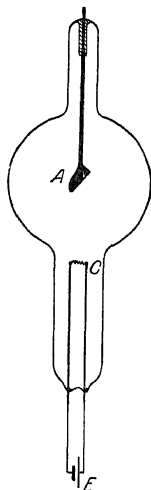


Fig. 9.—Diagrammatic representation of an X-ray tube. C, cathode; A anode (anticathode); E, battery for heating the cathode.

X-rays are produced when a cathode ray impinges on the glass wall of the tube, or on a specially fitted anticathode (fig. 9). Their penetrating power, also called their hardness, increases as the potential exciting the tube is raised. But different substances show different degrees of transparency. The higher the atomic weight, the greater the opacity—hence the possibility of obtaining radiographs of the bones, which contain a much greater proportion of metallic atoms than the surrounding flesh. The question, whether the rays are of

corpuscular nature, or waves similar to light, occupied physicists in vain for a long time. Interference and diffraction experiments (Walter and Pohl, 1908), the meaning of which we shall explain later (§ 1, p. 66), gave only one certain result, viz. that, if the rays are waves, their wave-length must be considerably shorter than that of visible or of ultraviolet light (the latter as shown photographically). By using a crystal as the diffracting apparatus, von Laue and his co-workers Friedrich and Knipping (1912) succeeded in deciding the question (p. 68): X-rays are light of very short wave-length. Properly speaking, therefore, they fall outside our present subject, which is the investigation of the elementary constituents of matter. But we shall often have to refer to them as powerful aids to the investigation of matter. It is very remarkable that by means of these very X-rays experiments have been carried out the results of which are completely opposed to the wave view, and have compelled us to interpret light in terms of corpuscles (§ 5, p. 77).

#### 4. Radiations from Radioactive Substances.

We have hitherto been occupied only with radiations produced artificially. As we know, there are also natural radiations, which are emitted by *radioactive substances*, the process involving spontaneous change of the atoms of these substances into other atoms (Becquerel, 1896; P. and M. Curie, 1898; Rutherford, Soddy, 1902). We distinguish three different kinds of radioactive radiations:

1.  **$\alpha$ -rays:** The deflection experiments show that in these we have to do with positively charged particles, which are much more difficult to deflect than cathode rays, and must therefore be of much greater mass than electrons. Their  $e/m$  ratio corresponds to that of a doubly ionized helium ion, i.e. to a  $\text{He}^{++}$  particle. That it is actually  $\text{He}^{++}$  (atomic weight 4) which is in question, and not, say, a singly charged particle of atomic weight 2, which would of course show the same value for the specific charge, is made very probable by the fact (Rutherford, Ramsay and Soddy, 1903) that radioactive substances develop helium, and has been proved in most convincing fashion in an experiment due to Rutherford and Royds (1909), who succeeded in capturing  $\alpha$ -particles in an evacuated vessel; when the gas composed of  $\alpha$ -particles is made luminous, the spectroscope reveals unmistakably the lines of helium.

2.  **$\beta$ -rays:** These, as deflection experiments show, undoubtedly consist of electrons, and differ from the cathode rays only by their higher velocity. While cathode rays can be produced with velocities

within a few hundredth parts of that of light, the velocity of  $\beta$ -rays differs from that of light by a few thousandth parts only.

More exact investigation has shown that two kinds of  $\beta$ -rays are emitted from a given radioactive element. One kind has a continuous "velocity spectrum", i.e. electrons occur with every possible velocity over an extensive range. The other kind has a discontinuous velocity spectrum, i.e. it consists of groups of electrons of definite velocity.

We shall see later that the latter sort can be regarded as a secondary effect of slighter importance; they do not originate at all in the nucleus, which is characteristic of the atom (p. 53), but in the external electronic system surrounding the nucleus. The nuclear  $\beta$ -rays proper, which have a continuous spectrum, present the theoretical physicist with a puzzling problem. For if particles of all possible energies actually leave the nucleus of the atom, the nucleus cannot remain as a unique structure of definite energy. But all experiments bear evidence that the product of the explosion process—the new atomic nucleus, poorer by an electron—is well-defined and unique and has therefore a definite energy. What happens then to the remainder of the energy? A direct proof of its survival (say as  $\gamma$ -rays) has not been found; we are faced here with a difficulty of principle (see p. 58).

3.  $\gamma$ -rays: These are found to be incapable of deflection in the electric or magnetic field. We have to do here with radiation of extremely short wave-length (ultra-X-radiation).

The fundamental law of radioactive transformation (v. Schweidler, 1905) states that the number of atoms disintegrating per unit time ( $-dN/dt$ ) is proportional to the number of atoms ( $N$ ) present at the moment;

$$-dN = \lambda N dt.$$

The factor of proportionality  $\lambda$  is called the *radioactive decay constant*; it is characteristic of the kind of atom. By integration we find

$$N = N_0 e^{-\lambda t},$$

where  $N_0$  denotes the number of atoms at the time  $t = 0$ . We can express  $\lambda$  in terms of the half-value period  $T$ , i.e. the time it takes for half of the atoms to be disintegrated. We have  $N_0 e^{-\lambda T} = \frac{1}{2} N_0$ ,

$$\text{or} \quad T = \frac{\log_e 2}{\lambda} = \frac{0.6931}{\lambda}.$$

For radium itself,  $T = 1590$  years; but there are radioactive elements

with extremely long half-value periods, as thorium with  $T = 1.8 \times 10^{10}$  years, and some with extremely short ones, like thorium  $B'$  with  $T = 10^{-9}$  sec.

The meaning of the law of transformation is, that every atom in a certain measure has the same explosion probability; clearly the law is a purely statistical one. This has been confirmed in two different ways. First, it has been found quite impossible by ordinary physical means (say high temperatures) to accelerate or retard the process of disintegration, or to affect it in any way. Secondly, it has been found possible to determine not only the mean number of particles emitted per second, but also the fluctuations about the mean; and it turns out that these obey the regular statistical laws (Appendix IV, p. 266). Radioactive disintegration is the prototype of an elementary process, which the ideas of classical physics are powerless to explain, but with which modern quantum theory is quite capable of dealing.

We shall now cite a few more experiments, which seem to point unambiguously to the corpuscular nature of the  $\alpha$ - and  $\beta$ -rays. We attach special weight to the fact that it seems simply impossible to understand these experiments from any other point of view than that we are actually dealing with discrete particles. In the next chapter, however, we shall discuss a series of experiments on these same rays which seem to indicate just as indubitably that the rays represent a wave process.

We begin with *scintillation phenomena* (Crookes, 1903), which have already been referred to at the end of the last chapter in connexion with methods of determining Avogadro's number. If a fluorescent screen is set up near a radioactive preparation, flashes of light are observed, now here, now there, on the material of the screen. Anyone who has a watch with a "self-luminous" dial can, with the help of a magnifying glass, convince himself of the presence of these flashes of light; the illuminating substance consists of a layer of radioactive material, varnished with zinc blende. When the radiation strikes the zinc blende, the spot is lit up. These phenomena compel us to assume that the  $\alpha$ -radiation consists of discrete particles, like a hail of shot, and that the fluorescent screen becomes illuminated wherever it is struck by such a particle.

A second vivid proof of the corpuscular nature of these rays is given by another method of counting the individual particles. The Geiger counter (1913; Geiger and Müller, 1928) consists essentially of a metal plate, placed opposite a metallic point (fig. 10); the whole is contained in an air-filled vessel. A potential difference is applied to the plate and point, as great as possible subject to the condition that in spite of the

action of the point no discharge passes. If an  $\alpha$ -particle now flies past, it will ionize the molecules of air which it strikes, and this minute alteration in the distribution of the field just suffices to produce a discharge between plate and point. After this discharge, the plate and point again become charged, and the process begins afresh. Every particle which flies past causes in this way a momentary discharge which can be recorded by suitable apparatus, or, with the help of an amplifier, heard in a telephone.

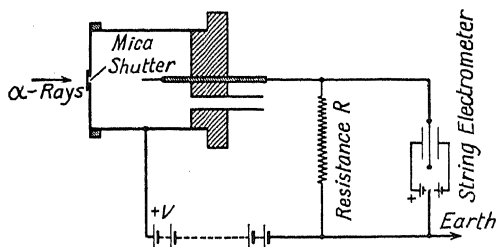


Fig. 10.—Diagrammatic representation of a Geiger counter for recording  $\alpha$ - and  $\beta$ -particles. The applied potential is just great enough to prevent spontaneous discharge between the point and the wall of the instrument. An ionizing particle flying in starts the discharge; a resistance  $R$  serves to limit the current and thus break off the discharge. The discharges are counted, e.g., by means of a recording arrangement (amplifier and telephone) inductively connected to the circuit; or, as in the figure, by the throws of a string electrometer.

For proving the corpuscular nature of the radiations from radioactive bodies, the most picturesque method consists in making the individual particles directly visible in the *Wilson cloud chamber* (1912) (fig. 11).

If we take pure water vapour in which there are no nuclei on which condensation could take place, in the form of dust particles and so on, and if by sudden *expansion* we cool it sufficiently for the vapour pressure of the water to be exceeded, then in the absence of nuclei of condensation the vapour cannot condense

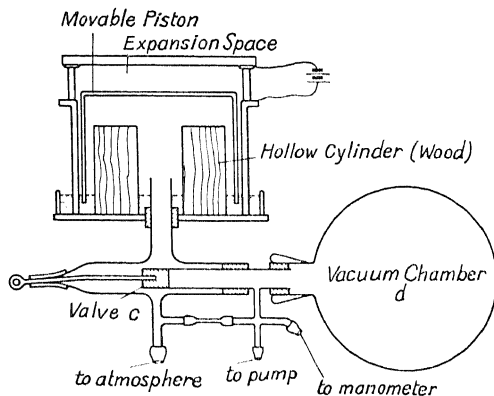


Fig. 11.—Diagrammatic representation of the Wilson cloud chamber. The movable piston is suddenly lowered by opening the valve  $c$  and so connecting the vacuum chamber  $d$  with the part of the apparatus beneath the piston.

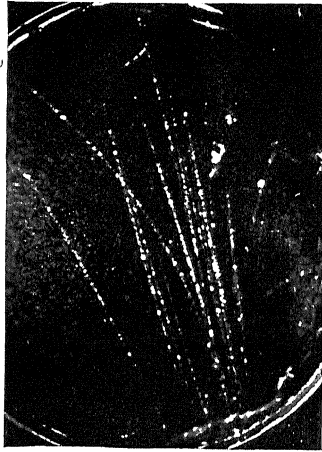
into drops; we obtain supersaturated water vapour. If an  $\alpha$ -particle shoots into the gas in this condition, it will ionize the molecules with which it collides. But the ions now act as nuclei of condensation, on which the neighbouring molecules of the supersaturated vapour are

deposited as droplets. The path of the  $\alpha$ -particle is thus made visible in the form of a series of minute water drops. An ingenious arrangement makes it possible by a single operation to expand the vapour beyond saturation point, to give the radioactive rays free access to the vapour, to expose the whole apparatus to a flash of light, and to take an instantaneous photograph. The plate shows a series of rectilinear or broken paths. The paths of the  $\alpha$ -particles can be distinguished from those of the  $\beta$ -particles by the greater ionizing action of the former (more intense condensation), and also by their greater penetrating power corresponding to their greater mass (the paths thicker, straighter, less broken). By collecting a sufficient amount of photographic material, we can make statistical statements about the mean length of the paths, the size of the breaks, &c.

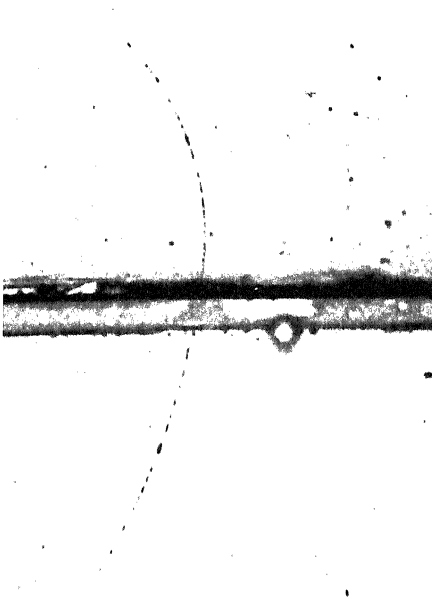
### 5. Prout's Hypothesis, Isotopy, the Proton.

Now that we have sufficiently convinced ourselves of the corpuscular character of the  $\alpha$ - and  $\beta$ -rays from a radioactive substance, we proceed to consider the question which was raised at the outset, concerning the building up of atoms and molecules from elementary constituents. So far back as the beginning of last century the hypothesis was advanced by Prout that all atoms are ultimately *made up of hydrogen atoms*. It fell into oblivion, however, when chemists became able to determine the atomic weights of the elements more exactly. If Prout's hypothesis were correct, the atomic weight of every atom would of course be a whole number of times that of the hydrogen atom. It is found, however, that for a whole series of atoms this is not the case; a glaring example is chlorine, the atomic weight of which, referred to that of hydrogen as unity, has the value 35.5.

The phenomena accompanying radioactivity were responsible for the conjecture that elements, though chemically absolutely pure, actually represent a mixture of different kinds of atoms, of the same structure indeed, but of different mass. These atoms, perfectly equivalent chemically, but of different mass, are called *isotopes*. To make what follows more easily understood, we should like briefly to remind the reader of the *periodic system* of the elements (Table I, p. 35). The chemical behaviour of a given element is to a large extent determined by the place which it occupies in the periodic table; thus the alkali metals, with their chemically similar behaviour, occupy the same vertical column, likewise the alkaline earths, the noble metals, the heavy metals, and finally the halogens and the inert or noble gases. The arrangement of the elements in this system was originally carried out on the



Ch. II, Fig. 14.—“Showers” of electrons and positrons liberated by cosmic rays  
(From *Proc. Roy. Soc. A*, 139, Plate 22.)



Ch. II, Fig. 12. A 63-million volt positron passing through a 6-mm. lead plate and emerging as a 23-million volt positron.

(*Handbuch der Physik*, Julius Springer.)





TABLE I.—PERIODIC TABLE OF THE ELEMENTS

	I	II	III	IV	V	VI	VII	VIII
1	1 H 1.0078							2 He 4.002
2	3 Li 6.940	4 Be 9.02	5 B 10.82	6 C 12.00	7 N 14.008	8 O 16.0000	9 F 19.000	10 Ne 20.183
3	11 Na 22.997	12 Mg 24.32	13 Al 26.97	14 Si 28.06	15 P 31.02	16 S 32.06	17 Cl 35.457	18 A ← 39.944
→	19 K 39.096	20 Ca 40.08	21 Sc 45.10	22 Ti 47.90	23 V 50.95	24 Cr 52.01	25 Mn 54.93	26 Fe 55.84
4	29 Cu 63.57	30 Zn 65.38	31 Ga 69.72	32 Ge 72.60	33 As 74.91	34 Se 78.96	35 Br 79.916	27 Co ← → 28 Ni 58.69
	37 Rb 85.44	38 Sr 87.63	39 Y 88.92	40 Zr 91.22	41 Nb 93.3	42 Mo 96.0	43 Ma	44 Ru 101.7
5	47 Ag 107.880	48 Cd 112.41	49 In 114.76	50 Sn 118.70	51 Sb 121.76	52 Te 127.61	53 I 126.92	45 Rh 102.91
	55 Cs 132.91	56 Ba 137.36	57 La 138.92	72 Hf 178.6	73 Ta 181.4	74 W 184.0	75 Re 186.31	46 Pd 106.7
6	79 Au 197.2	80 Hg 200.61	81 Tl 204.39	82 Pb 207.22	83 Bi 209.00	84 Po (210.0)	85 —	76 Os 191.5
7	87 —	88 Ra 225.97	89 Ac (227)	90 Th 232.12	91 Pa (231)	92 U 238.14		77 Ir 193.1
								78 Pt 195.23
								86 Rn 222

THE RARE EARTHS (to be inserted between 57 La and 72 Hf)

58 Ce 140.13	59 Pr 140.92	60 Nd 144.27	61 —	62 Sm 150.43	63 Eu 152.0	64 Gd 157.3
65 Tb 159.2	66 Dy 162.46	67 Ho 163.5	68 Er 167.64	69 Tm 169.4	70 Yb 173.04	71 Lu 175.0

The numbers in front of the symbols of the elements denote the atomic numbers; the numbers underneath are the atomic weights. The latter are taken, with a few modifications, from the Report of the International Commission on Atomic Weights for 1932. The double arrow ← → indicates the places where the order of atomic weights and that of atomic numbers do not agree.

basis of their atomic weight; to-day we arrange them according to their "atomic number", as we shall see below.

At the end of the periodic system stand the radioactive elements; these too can be arranged according to their chemical behaviour. The method, first applied by the Curies, for separating them and determining their chemical properties, consists in bringing the radioactive substance with other bodies into solution, and applying various precipitating agents. It is then tested whether the activity, recognizable by the radiating power, is in the precipitate or in the residual solution, or is divided between the two. Each portion is treated in a similar way, until one part of the products is free from radiation, and the other part shows an exponential decay with characteristic half-value period (§ 4, p. 31). Radium itself was isolated thus by the Curies (1898), starting from the mineral pitchblende; barium was left along with the radium to begin with, and could not be separated from it until the end. Radium therefore belongs to the group of alkaline earth metals, and so must be placed in the periodic table in the column under calcium and barium; radon Rn (or Ra emanation), on the other hand, behaves like a noble gas and so ranges under He, Ne, &c. The short-lived succeeding disintegration product RaA ( $T = 3$  min.) is found to be chemically analogous to tellurium; the next, RaB ( $T = 27$  min.), goes with lead, and so on.

From these and similar results the following important *law of radioactive transformation* has been deduced: emission of an  $\alpha$ -particle (loss of charge  $+2e$ ) shifts the residual atom two places to the left in the periodic table, i.e. in the direction towards lower H-valencies; thus radium on giving up an  $\alpha$ -particle is transformed into radon (Ra emanation). On the other hand, escape of a  $\beta$ -particle (loss of charge  $-e$ ) displaces the atom one place to the right (Russel, Soddy, Fajans, 1913). If we go through the three radioactive series according to this rule (fig. 13), we find that many places in the table are multiply occupied. To cite a particularly conspicuous example: into the place in the periodic table occupied by ordinary lead, there also come, according to the law of radioactive change, the end-products of the three radioactive series, viz. RaG, AcD and ThD, all three of which have undoubtedly the character of lead. Besides these, RaD, AcB, ThB and RaB also fall into the place mentioned. All these elements, in spite of their like chemical behaviour, have different masses, because in an  $\alpha$ -emission the atom gives up the mass 4 of helium, while in a  $\beta$ -disintegration, on account of the small mass of the electron, the mass of the atom remains practically unchanged. The three elements of the radium

series mentioned above, RaB, RaD and RaG, must therefore have masses decreasing successively by 4.

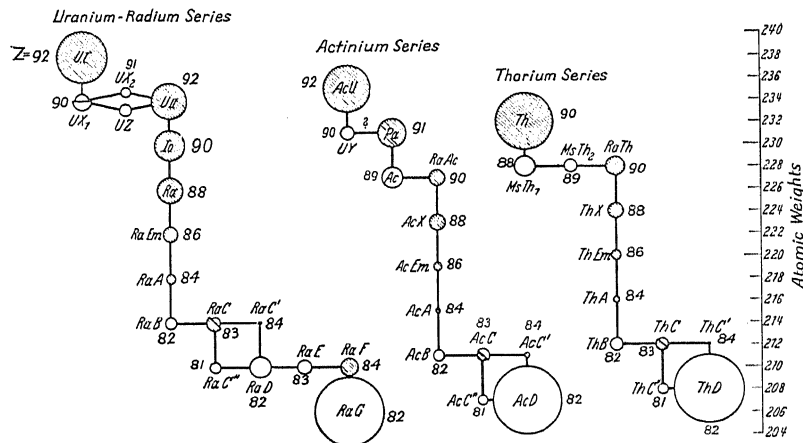


Fig. 13.—The radioactive transformation series. The changes indicated by vertical lines correspond to  $\alpha$ -emission, by horizontal lines to  $\beta$ -emission; in the former the atomic weight falls by 4, the atomic number by 2; in the latter the atomic weight remains approximately constant, and the atomic number increases by 1. Shaded circles indicate  $\alpha$ -rays; circles without shading,  $\beta$ -rays. The size of each circle corresponds to the half-value period.

It may also happen, we may add (e.g. in every  $\beta$ -disintegration), that two chemically different elements have the same atomic weight; these are called *isobars*.

After the existence of isotopes had been demonstrated in this way for radioactive substances, J. J. Thomson (1913), by deflection experiments on canal rays, succeeded in proving that isotopes occur even among ordinary elements, for instance neon. Ordinary chlorine, whose atomic weight is given by chemists as 35.5, consists of one kind of chlorine of weight 35.0, and another of weight 37.0. Thomson's investigations reached their highest development in the mass-spectrograph constructed by Aston (1919), which we have already mentioned (§ 2, p. 29), and which at present represents the most exact method for the determination of atomic weights (Aston, Dempster (1918), Bainbridge).

By magnetic separation of canal rays of lithium, it has even been found possible (Oliphant, Shire and Crowther, 1934) to obtain visibly separated deposits of the two lithium isotopes (6 and 7), which can be employed for other experiments.

Recently (1933) G. Hertz has succeeded in separating isotopes by a

mechanical method. By means of a suitably devised circulation process, he caused the mixture of isotopes to diffuse several times through a system of clay cylinders. Since the lighter components of the mixture diffuse more rapidly than the heavier, we obtain in this way, as the final products of the circulation process, two mixtures, one of which is richer in the lighter components, the other in the heavier. This method has been applied successfully for the separation of the heavy isotope of hydrogen, of which we have to speak later (§ 5, p. 59).

The result of the detailed investigations on the occurrence of isotopes may be stated in the following form. Every element (in the chemical sense) has either a whole number for its atomic weight, or consists of a mixture of different kinds of atoms whose atomic weights are whole numbers. For practical reasons atomic weights have all along been referred, not to hydrogen equal to 1, but to oxygen equal to 16. This has proved fortunate as an aid to clearness, for the integral character of the isotopic weights stands out much more obviously than with the choice of  $H = 1$ . Hydrogen indeed, for  $O = 16$ , gets a value ( $H = 1.0081$ ) differing decidedly from 1, but the atomic weights of all other pure isotopes approximate closely to whole numbers.

To understand how this comes about, we must anticipate so far as to note that the masses of atoms are almost entirely concentrated in the *atomic nuclei*. Every atom consists of a nucleus, which is surrounded by a cloud of electrons (§ 3, p. 51). This nucleus is the essential part of the atom. The result of Aston's investigations must be referred to the nucleus: every nucleus consists of a whole number of hydrogen nuclei or *protons*; its mass (except for trifling divergencies) is an integral multiple of the mass of the proton.

These divergencies, which in spite of their slightness are of the highest importance, will be discussed later (§ 4, p. 55).

Here we may once again summarize the ideas which for something like two decades dominated physics.

There are two primitive atoms, the atoms of electricity, the negative electron and the positive proton; they have equal and opposite charges, but (very remarkably) quite different masses (in the ratio 1:1840; see § 1, p. 27). From these all matter is built up, and that, as we shall explain more fully in next chapter, in two stages: there is first formed, from protons with some cementing electrons, the very small and compact nucleus; then this is surrounded by a cloud of electrons of relatively loose structure.

But this simple and homogeneous picture has, in the light of a series of new discoveries, turned out to be incorrect, as we must now explain.

TABLE II.—TABLE OF ISOTOPES

The isotopes are arranged in each case where known in the order of frequency of occurrence; the radioactive isotopes are indicated by an asterisk. Radioactive isotopes produced artificially are not included.

Element	Z	Isotopes	Element	Z	Isotopes
H	1	1, 2, 3	Sn	50	120, 118, 116, 119, 117, 124, 122, 112, 114, 115
He	2	4, 3	Sb	51	121, 123
Li	3	7, 6	Te	52	130, 128, 126, 125, 124, 122, 123, 127?
Be	4	9	I	53	127
B	5	11, 10	Xe	54	129, 132, 131, 134, 136, 130, 128, 124, 126
C	6	12, 13	Cs	55	133
N	7	14, 15	Ba	56	138, 135, 136, 137, 134, 132, 130
O	8	16, 18, 17	La	57	139
F	9	19	Ce	58	140, 142, 138, 136
Ne	10	20, 22, 21	Pr	59	141
Na	11	23	Nd	60	146, 144, 142, 145, 143
Mg	12	24, 25, 26	Sm	62	144, 147, 148, 149, 150, 152, 154
Al	13	27	Eu	63	151, 153
Si	14	28, 29, 30	Gd	64	155, 156, 157, 158, 160
P	15	31	Tb	65	159
S	16	32, 34, 33	Dy	66	161, 162, 163, 164
Cl	17	35, 37	Ho	67	165
A	18	40, 36, 38	Er	68	166, 167, 168, 170
K	19	39, 40*, 41	Tm	69	169
Ca	20	40, 44, 42, 43	Yb	70	171, 172, 173, 174, 176
Sc	21	45	Lu	71	175
Ti	22	48, 50, 46, 47, 49	Hf	72	176, 178, 180, 177, 179
V	23	51	Ta	73	181
Cr	24	52, 53, 50, 54	W	74	184, 186, 182, 183
Mn	25	55	Re	75	187, 185
Fe	26	56, 54, 57, 58	Os	76	192, 190, 189, 188, 186, 187
Co	27	59	Ir	77	193, 191
Ni	28	58, 60, 62, 61, 64	Pt	78	196, 195, 194, 192, 198
Cu	29	63, 65	Au	79	197
Zn	30	64, 66, 68, 67, 70	Hg	80	202, 200, 199, 201, 198, 204, 196
Ga	31	69, 71	Tl	81	205, 203, 207*, 208*, 210*
Ge	32	74, 72, 70, 73, 76	Pb	82	208, 206, 207, 204, 203?, 205?, 209?, 210*, 211*, 212*, 214*
As	33	75	Bi	83	209, 210*, 211*, 212*, 214*
Se	34	80, 78, 76, 82, 77, 74	Po	84	210*, 211*, 212*, 214*, 215*, 216*, 218*
Br	35	79, 81	Rn	86	222*, 210*, 220*
Kr	36	84, 86, 82, 83, 80, 78	Ra	88	226*, 223*, 224*, 228*
Rb	37	85, 87*	Ac	89	227*, 228*
Sr	38	88, 86, 87	Th	90	232*, 227*, 228*, 230*, 234*
Y	39	89	Pa	91	231*, 234*
Zr	40	90, 94, 92, 96, 91	U	92	238*, 235*, 234*
Cb	41	93			
Mo	42	98, 96, 95, 92, 94, 100, 97			
Ru	44	102, 101, 104, 100, 99, 96, 98?			
Rh	45	103			
Pd	46	102, 104, 105, 106, 108, 110			
Ag	47	107, 109			
Cd	48	114, 112, 110, 113, 111, 116, 106, 108			
In	49	115, 113			

## 6. The Neutron.

The helium nucleus of mass 4 and charge 2 will, according to the ideas just developed, consist of 4 protons and 2 electrons; and similarly in higher nuclei some electrons will partly compensate the sum of the proton charges. The question now arises: why is there always a surplus positive charge? May not a nucleus exist with equal numbers of protons and electrons, e.g., as the simplest form, a neutron, consisting of a proton and an electron?

Reflections of this sort have been dwelt upon by many writers, in the course of speculations on the nature of the nucleus and of radioactive disintegration. The actual discovery of neutrons was due to purely experimental results, referring in fact to artificial transformations of atoms. Thus it was observed by Bothe and Becker (1930), when light elements like lithium and beryllium were bombarded by  $\alpha$ -rays, that  $\gamma$ -rays were emitted. Then Irène Curie, daughter of the discoverers of radium, and her husband Joliot found (1932) that the radiation from the bombarded beryllium, when passed through substances containing hydrogen, such as paraffin, expelled protons, which  $\gamma$ -rays never do; a new kind of radiation must therefore also be present. From an investigation with an ionisation chamber it was inferred by Chadwick (1932) that the new radiation consists of uncharged, heavy particles. This was confirmed by Feather (1932) with help of the Wilson Chamber. The track of the particle coming from the beryllium remains invisible, showing that it does not ionize the molecules of air, while of course we see the tracks of the nuclei on which it impinges. By comparison of the ionisation in nitrogen and hydrogen, Chadwick was able to estimate the mass of the neutron, finding it approximately equal to that of a proton.

With regard to their penetrating power, neutrons behave quite differently from any other kind of radiation, whether light waves or charged particles. With the latter, the process of absorption essentially consists in their giving up energy to the outer electrons of the atoms; as the number of these runs roughly parallel with the mass, so also the absorbing power of different substances runs roughly parallel with their mass. The neutrons, however, pay no attention whatever to the outer electrons—only collisions with a nucleus stop them; since the size of the nucleus varies very little from one substance to another (p. 183), it is the number of nuclei per unit volume which now counts. But 1 gm. of hydrogen contains the same number of nuclei as 16 gm. of oxygen; the absorbing power per gramme is therefore some 16 times as great for hydrogen as for oxygen. The Wilson photo-

graphs show, moreover, that in some collisions of neutrons and nuclei the neutron flies on (it is supposed that its diameter is of the same order of magnitude as that of a nucleus), while in others it is captured by the nucleus, and causes this to explode, with expulsion of other particles (§ 6, p. 61). We thus come to the question of what part the neutrons take in the structure of nuclei; we shall deal with this later (p. 57), and only remark here that of course it is now possible to regard all nuclei, of whatever (positive) charge, as made up of protons and neutrons alone, without making any use of electrons.

The electron would then only have a place in the nucleus as a constituent of the neutron. But the question whether even this supposition is valid, viz. that the neutron = proton + electron, is by no means easy to answer. For meanwhile the position has become further complicated by the discovery of the positive electron, or *positron*, of which we shall speak presently. It would now even be possible to regard the neutron and positron, along with the electron, as the elementary constituents of matter, and to put the proton equal to neutron + positron. The question—if it has a meaning at all—is one of stability, i.e. of energy balance; for energy is required to separate two particles, and we can therefore decide the question by careful energy measurements; we shall return to this later (Chap. III, § 4, p. 55).

## 7. Cosmic Rays. Positrons.

There is no question which has presented (and still presents) greater difficulties to theoretical physicists than the difference in the masses of the positive and negative atoms of electricity. Every theory proposed up to the present is symmetrical in the sign of the charge. It would therefore long ago have been reasonable to conjecture that there are also particles of both kinds, heavy and light, with opposite charges to those known: positive electrons and negative protons. To advance this as a proposition was reserved, however, for the latest form of quantum mechanics (1928); the existence of positive electrons and the conditions for their occurrence and disappearance follow from Dirac's relativistic wave mechanics of the electron (§ 8, p. 169).

The actual discovery of positive electrons or “positrons” was achieved independently of this, in connexion with observations on the so-called “cosmic rays”. These are of great interest for their own sake, so something may be said here about their discovery and investigation.

As has already been mentioned (§ 1, p. 24), gases in their natural state are poor conductors of electricity. But they can be made conducting in

various ways, not only, as explained above, by applying potentials at reduced pressure, but also by irradiation with all kinds of corpuscular rays or (short-wave) electromagnetic rays. These rays separate electrons from the atoms or molecules. These electrons, as well as the residues, the ions, are set in motion by an electric field. To demonstrate the presence of radiations, and investigate their properties, ionization chambers—small, gas-filled vessels, with electrodes—are used. In measuring the strength of the radiation by the ionic current, we must be certain that this current vanishes when there is no radiation. It is found, however, that the current is never completely absent. The residual ionization had been attributed to a weak terrestrial radioactive radiation. Thick lead shielding, however, does not suppress it entirely; a weak ionization always remains (Rutherford, McLennan, 1903), which must arise from a radiation of much higher penetrating power than any known  $\gamma$ -rays. It was observed by Gockel (1909) that this radiation did not diminish when the apparatus was taken to a height above the ground, in the way it ought to do if it originated in radioactive sources in the earth. Hess then (1912) showed by balloon ascents at heights up to 5 km. that the intensity of the radiation even increases with height; he found also that it is just as great by night as by day (and so cannot come from the sun). The rays appear therefore to come from interstellar space; they are called *cosmic rays*. The experiments were carried to greater and greater heights by Kohlhörster (1914), later by Millikan, Tizzard, and Regener; with recording balloons Regener (1935) reached as high as 30 km. The existence of the rays was also demonstrated at the bottom of deep lakes, up to 500 metres under water (Millikan, Regener, 1928).

More exact researches were carried out with the help of the Wilson cloud chamber (first by Skobelzyn, 1929) and a large magnetic field; the tracks of the particles were seen as circles of slight curvature; their velocity or energy was deduced by measurement of the radius. It was thus found that the rays consist of extremely swift particles; among them it is no rarity to find electrons with velocities to produce which would require a potential of 100 or 1000 million volts.

What we observe here at the earth's surface is certainly not the original cosmic radiation, but a mixture of that with secondary radiation, consisting of electrons ejected from air molecules, or appearing in the process of pair production referred to below. Some light appears to be thrown on the question of the nature of the primary radiation by the fact that its intensity depends on the geographical latitude (Clay); it is weaker at the equator than at the poles. It must therefore be



a case of electrically charged particles, which are deflected in the magnetic field of the earth; for a swarm of electrons, which strikes a magnet, concentrates chiefly round the poles—a fact which of course has long been known, from the explanation of the aurora by Störmer (since 1903). The cosmic rays seem to come pretty uniformly from all parts of the sky. It is too early to speculate on their origin.

Anderson was the first to remark that frequently two kinds of track appeared on a photograph, which were of opposite curvature, though otherwise both looked like electron tracks (fig. 14, Plate II). If then we do not believe in positive electrons, we must assume that the one track is traversed in the opposite sense from the other. This possibility was excluded as follows. A lead plate was set up in the chamber, and tracks were found due to particles which had passed through the plate. But such a particle must necessarily have a smaller velocity, and therefore a greater curvature of path, after it has passed through than before; so that the sense in which the tracks are described can be determined (fig. 12, Plate II). The existence of positrons was thus proved. Blackett and Occhialini (1933) found that frequently a whole shower of particles is ejected from the wall of the vessel, their tracks showing, some positive, some negative curvature. Since we must attribute their common origin to an atomic explosion, we have here new evidence for positive electrons. Recently their occurrence has been demonstrated in various processes of nuclear disintegration; we shall return to this later (§ 6, p. 61). It has been found further, that the primary cosmic radiation consists predominatingly of positive particles, presumably positrons; this is inferred from the fact that in equatorial latitudes on the earth more rays come from the west than from the east.

From the theoretical point of view the discovery of the positron is of the highest importance; for it confirms the theory of Dirac already mentioned, which on the intuitional side amounts to this, that neither electrons nor positrons are immutable elementary particles, but that by coming into collision they may annihilate each other, with the emission of energy in the form of light waves; and conversely, that a light wave of high energy can in certain circumstances become the source of a “pair” (electron + positron). There is experimental evidence for both processes. The fact that in our actual world negative electrons preponderate, is not inconsistent with the theory.

## CHAPTER III

### The Nuclear Atom

#### 1. Lorentz's Electron Theory.

The method of physical science in the investigation of the structure of matter has in all times been based on the following principle: laws established for "macroscopic bodies", i.e. bodies of ordinary size, are applied tentatively to elementary particles; if some disagreement is then found, an alteration of the laws is taken in hand. In this way, advance essentially depends upon the closest co-operation of observation and theoretical interpretation.

Thus in the preceding pages we have made continual use of the laws of interaction of charged particles, and of their response to the influence of external fields and light rays. The most important result, that the mass of an atom is almost wholly concentrated in a very small nucleus, while its volume, and its physical and chemical properties, are determined by a comparatively loose surrounding structure of electrons, we have assumed without proof. This proof we must now supply. First, however, we must look a little more deeply into the laws of the electromagnetic field.

As we know, these are formulated in Maxwell's equations (1855) (App. VI, p. 272). As originally stated, these equations represent the changes in space and time of the electric field strength  $\mathbf{E}$ , and the magnetic  $\mathbf{H}$ , in material bodies, and therefore contain certain constants which are characteristic of those bodies (in the simplest case they are the dielectric constant and the magnetic permeability).

The atomic theory of electricity simplifies the field equations, by considering only fields in a vacuum; the units can then be chosen so that the field equations contain no material constants (or at most, the velocity of light  $c = 3 \cdot 10^{10}$  cm./sec., which is retained so that we may be able to use the ordinary units of kinematics, centimetres and seconds). This is the standpoint of Lorentz's electron theory (1897), which dominated physics about the turn of the century. At that time we were

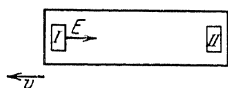
acquainted with free electrons and were aware that they are constituents of atoms; about the positive atoms of electricity nothing was known. The electrons were pictured as very small charged bodies, which generated the field in free space, and conversely were acted on by forces due to the field. These forces determined the motion of the electrons in accordance with Newton's law: mass into acceleration equals force. It was therefore necessary to ascribe to the electrons a mass  $m$  besides their charge  $e$ , and we have seen that the fraction  $e/m$  can be determined by deflection experiments. On the basis of these assumptions the motion of electrons in various fields could be calculated; it was thus found that the size of the electron (its radius, if it was pictured as a sphere) plays a part, and that for the following reason. When there is acceleration, the field of the charges in the electron is modified, and this changed field passes over the electronic sphere with the velocity of light, and exercises forces on it. For small accelerations, the forces to a first approximation are proportional to the acceleration and to  $e^2/a$ , where  $e$  is the charge, and  $a$  the radius (with a numerical factor which depends on the distribution of the charge); in other words, the effect is the same as if the mass were increased by a part proportional to  $e^2/a$  (J. J. Thomson, 1882; Heaviside, Searle, 1885). This fact suggested to some physicists the idea that the electron possesses no "ordinary" mass at all, but only "electromagnetic" mass; and the smallness of the mass of the electron seemed to confirm this. Further, the electromagnetic mass was found to depend on the velocity, and the first observations of this effect by Kaufmann (1906), were regarded as a brilliant triumph for the theory. Hasenöhl (1904) derived the laws of motion of an empty box with reflecting walls filled with electromagnetic radiation, and found that this radiation behaved as having a mass  $E/c^2$ , where  $E$  is the total electromagnetic energy. This was the first indication of the general law considered in the following section.

The calculations of the electromagnetic mass of the electron rested on the assumption (Abraham, 1903) that the electron is rigid, retaining its form throughout the motion; which implies the assumption of infinitely great internal forces of non-electromagnetic origin. If the assumption is dropped, we obtain not only other numerical factors, but also other functions of the velocity. Besides, the assumption of absolute rigidity is quite incompatible with the theory of relativity (on account of the Lorentz contraction; see Appendix V, p. 269); if instead of this we postulate invariability of form in the reference system in which the electron is instantaneously at rest, we obtain a

formula which agrees, to a numerical factor, with that stated on p. 27 (Lorentz's electron, 1909). The numerical factor, however, must to a large extent remain arbitrary, since it depends on the distribution of the charge, as to which we know nothing. Worse still, even in this relativistic form very great internal forces of cohesion must be assumed, to keep the parts of the electron together, these having charges of the same sign and therefore repelling each other; and this leads to a contradiction of a fundamental theorem of the mechanics of radiating systems, of which we have now to speak.

## 2. The Theorem of the Inertia of Energy. Unitary Field Theory.

The theorem of the inertia of energy in its full generality was first stated by Einstein (1905). It asserts that any energy  $E$  possesses a mass  $m$ , in accordance with the equation  $E = mc^2$ . Of this relation we shall later continually make use. The relation is very far from being a mere theoretical subtlety; the phenomena connected with it have as a matter of fact the character of large scale effects. Think of a closed box, in the sides of which are fitted two exactly similar instruments (I and II), which are so constructed that they can



send out a momentary light signal in a definite direction, or completely absorb an incoming light signal (fig. 1). Now let the instrument I, at a



Fig. 1.—Illustration of Einstein's ideal experiment to prove the relation  $E = mc^2$ . The transmitter I radiates to the receiver II a definite quantity of energy  $E$ ; in consequence of this the whole box undergoes a recoil.

definite moment, send out a light signal in the direction of the instrument II. During this process of emission, the instrument I, and with it the whole box, experiences a recoil. The occurrence of this recoil is due to the radiation pressure. The latter phenomenon was observed experimentally by Lebedew (1901), in good agreement with theory; it was investigated later by Nichols and Hull (1903) and others, and finally very exactly by Gerlach and his collaborators (1923). In consequence of the recoil, during the whole time which the light takes to go from I to II the box will move in the opposite direction, and will not come to rest until the light strikes the instrument II, and the radiation pressure on it again brings the box to a standstill. Now interchange the two instruments; this alters nothing, since of course the two were given the same mass. Then let the instrument II send out towards I the same quantity of light as was sent before by I, whereby the whole box is again displaced to the left the same distance as before. Now interchange the instru-

ments once more. By proceeding in this way, it would be possible to displace the box any distance, without any change taking place in the interior of the box or in its neighbourhood—a conclusion obviously at variance with a fundamental property of the centre of inertia.

The contradiction disappears at once, however, if we take account of Einstein's thesis of the equivalence of energy and mass. By the emission of the light signal at the first step of our ideal experiment, the first instrument gives out a definite quantity of energy  $E$ ; its content of energy, and accordingly its mass also, become smaller. Similarly, the energy content, and therefore the mass, of the instrument II is increased on its absorption of the light signal, so that the mass of II is now greater than that of I; to interchange I and II without shifting the centre of inertia therefore requires that the whole box should be displaced a definite distance to the right. If we examine now what the relation between mass and energy must be, in order that the displacement of the box, due to radiation pressure, should be exactly compensated by interchanging the two instruments, a short calculation (Appendix VII, p. 274) leads to Einstein's formula, as quoted above. The most important field of application of Einstein's law is the disintegration of nuclei; we shall see later (p. 62) that this gives an experimental proof of it.

We return now to the problem of electromagnetic mass (p. 45). According to Einstein's theorem, the simplest way of obtaining this mass for small velocities should be, to calculate the internal electric energy of the charge collected in the electron; this energy is in fact proportional to  $e^2/a$ , but the numerical factor is in all circumstances different from the one we get by calculating the force of reaction of the electron's own field. The contradiction, as we have already seen, is due to the presence of forces of cohesion, which also should make a contribution to the energy (though this is difficult to reconcile with the postulated rigidity).

In consequence of these unsatisfactory results, the tempting idea of electromagnetic mass has gradually been given up. It was found that the theory of relativity suggested in a purely formal way a law of dependence of mass on velocity for every body—a law which has been brilliantly confirmed by experiment (Bucherer, 1909; Neumann-Schäfer, 1914; Guije-Ratnowski-Lavanchy, 1921). As the quantum theory developed, physicists became sceptical about definite models of the elementary particles. They preferred therefore to think of the electron, in regard to all external actions, as a charged point mass, without troubling further about its internal structure. But even this point of view has

its difficulties. In the first place it means, not a solution, but a shifting of the problem. For the electron's proper energy, which is proportional to  $e^2/a$ , becomes infinite when  $a$  is put equal to 0. We must therefore eliminate the energy from all physical laws, a procedure which, under both classical methods and those of the quantum theory, leads to very artificial constructions, which all contradict the theorem of the inertia of energy; for the field energy even of a point charge is of course, according to the Faraday-Maxwell view, by no means concentrated in the electron, but must be thought of as distributed throughout space.

In the second place, the radius  $a$  of the electron has an actual physical significance; by  $a$  we understand that length which satisfies Einstein's relation  $a \cdot e^2/a = mc^2$ , where  $a$  is a numerical factor of the order of magnitude 1. We therefore have

$$a = a \cdot \frac{e^2}{mc^2} = a \cdot 2.817 \times 10^{-13} \text{ cm.}$$

The simplest phenomenon in which this quantity occurs is the *scattering of light* (or other electromagnetic radiation) by atoms. According to Maxwell's electromagnetic theory, light (as also X-rays) consists of a periodically variable, electromagnetic alternating field. If the light wave strikes a charged particle which can move freely, the latter is set vibrating, and that the more strongly, the lighter the particle is. If the particle is bound to other particles, its induced oscillation will have a greater amplitude the closer the frequency of the incident light is to its proper frequency. Hence, on the one hand, electrons, on account of their trifling mass, will vibrate in sympathy with the light much more strongly than protons, say, or still heavier particles; on the other hand, with visible light, to a very large extent only those particles which occupy places near the surface of the atom, and are therefore relatively loosely bound, will be excited to vibration; while for excitation of the bound electrons farther inside the atom, X-rays will be needed.

Now, as we know, a vibrating charged particle acts like an antenna—it sends out electromagnetic vibrations of its own in the form of spherical waves; and, in fact, the energy drawn per unit time from the primary ray and converted into scattered radiation by a (free or loosely bound) electron is given (Appendix VIII, p. 275) by

$$I = \frac{8\pi}{3} \left( \frac{e^2}{mc^2} \right)^2 I_0,$$

where  $I_0$  is the intensity of the primary radiation. Since  $I_0$  is defined as energy per square centimetre, but  $I$  denotes the whole scattered energy, the quotient  $I/I_0$  must have the dimensions of an area. We can put it equal to an "effective cross-section" of the electron  $\pi a^2$  and have then

$$\frac{I}{I_0} = \frac{8\pi}{3} \left( \frac{e^2}{mc^2} \right)^2 = \pi a^2,$$

so that

$$a = \sqrt{\frac{8}{3}} \frac{e^2}{mc^2}.$$

This quantity is the "radius" of the electron if  $a$  is put equal to  $\sqrt{(8/3)}$ .

The remarkable thing about this method is that it makes no use of any hypothetical extrapolation of electrostatics to the *interior* of the electron, but works with a point electron.

Quantum theory calculations of the scattering of light at free electrons, taking account of the theory of relativity (Klein-Nishina, 1929), also all lead to effective cross-sections, which depend on functions of this radius  $a$ .

Exactly analogous results hold for the scattering of electrons at nuclei (or other electrons), with which we shall deal in detail in the next section (p. 53). Here we shall merely put the question: how near can two electrons approach each other, if the velocity of each at a great distance is  $v$ ? Clearly the most favourable case is that in which they move towards each other in the same straight line. In this case, at the moment of closest approach, the velocity is zero, and therefore also the kinetic energy; the energy is therefore the sum of the rest energy  $2m_0c^2$  and the potential energy  $e^2/r$ . By the principle of the conservation of energy, this must be equal to the original energy,

which by the theory of relativity is  $2mc^2 = 2m_0c^2 / \sqrt{1 - \frac{v^2}{c^2}}$ ; thus we have

$$2m_0c^2 + \frac{e^2}{r} = 2m_0c^2 / \sqrt{1 - \frac{v^2}{c^2}},$$

whence

$$r = \frac{e^2}{2m_0c^2} \left/ \left( \frac{1}{\sqrt{1 - \frac{v^2}{c^2}}} - 1 \right) \right.$$

We see that here also the effective radius is proportional to  $e^2/mc^2$ .

It might be held that the electronic radius here is a mere artificial magnitude, obtainable by calculation from the constants  $e$ ,  $m$  and  $c$ ;

still, the newer developments show that all our theories appear to fail at lengths of this order of magnitude. Surely then we must assume that the radius of the electron has an actual meaning. The state of affairs requires a modification of electrodynamics of such a sort that the existence of the electronic radius becomes intelligible without bringing in the mechanically defined mass  $m$ . Attempts in this direction have not been lacking; the works of Mie (1912) in particular deserve special mention. The present writer (1933) has pointed out a way which leads to the required end without going beyond the limits of classical methods, and which can also be accommodated to the quantum theory.

The leading idea of this *unitary theory of field and matter* is, that there is a natural unit for the field strengths  $E$  and  $H$ , an absolute field strength  $b$ , which at the same time represents a sort of upper limit for all fields (like the velocity of light  $c$  for velocities). Formally, the new field equations are identical with those of Maxwell for an isotropic medium with a dielectric constant  $\epsilon$  and a permeability  $\mu$ , except that  $\epsilon$  and  $\mu$  are not constant, but functions of the field itself, viz. (Appendix VI, p. 272).

$$\mu = \frac{1}{\epsilon} = \sqrt{1 + \frac{1}{b^2} (H^2 - E^2)};$$

their product is 1, as in Maxwell's equations for free space. It is then found that (positive or negative) point charges exist (as singularities of the field), but that the corresponding field energy does not become infinite. From  $e$  and  $b$  we can calculate an "electronic radius"  $a$  by the formula

$$a = \sqrt{\frac{e}{b}};$$

the energy is then given by

$$E = 1.2361 \frac{e^2}{a},$$

with a numerical factor free from all arbitrariness. Such a charge moves in external fields, which do not change to any extent within distances of the order of magnitude of  $a$ , exactly like a Lorentzian electron with purely electromagnetic mass; and Einstein's theorem

$$E = mc^2$$

is rigorously valid (Born and Infeld, 1934). For short-waved fields (wave-length of order  $a$ ), however, deviations occur from the ordinary laws of motion of the electron. Similar considerations can be applied



in the quantum theory also. For every isolated electrodynamic system there exists a representative point which moves according to the laws of the quantum mechanics of particles. The mass is of course taken as equal to the energy, divided by  $c^2$ . There exists a finite angular momentum of the field the laws of motion of which are closely connected with those of the spinning electron, of which we have to speak later (see § 1, p. 136, and § 8, p. 169). The actual calculation of the smallest possible masses, it is true, has not yet been effected. Hence it is not yet settled whether this theory means a forward step towards understanding the elementary particles. In any case it abolishes the dualism of field and particles in a manner free from contradictions.

### 3. Investigation of Atomic Structure by Scattering Experiments.

The most important method for investigating atomic structure consists in causing radiation of some kind to fall on the atom, and observing how this is altered by the atoms. The alteration consists in a weakening (absorption) of the radiation which passes on undeviated, and the corresponding production of diffracted or scattered radiation. The case of the scattering of light at a free or weakly bound electron has already been considered (§ 2, p. 48). The process can be used for the purpose of counting the loosely bound electrons in an atom. If their number is  $n$ , and  $N$  is the number of atoms per unit volume, then the energy lost by the primary radiation per centimetre of its path, the absorption constant, is given (J. J. Thomson, 1906) by

$$\frac{nNI}{I_0} = \frac{8\pi}{3} \left( \frac{e^2}{mc^2} \right)^2 nN.$$

Since  $N$  is known from kinetic theory data,  $n$  can be determined from absorption measurements on X-rays; by taking these sufficiently hard, all the electrons, at least in the lighter atoms, can be regarded as practically free, so that  $n$  denotes the whole number of electrons, and so also the nuclear charge in terms of the elementary charge as unit.

The measurements showed, for all lighter atoms which could be investigated, that  $n$  is approximately equal to half the atomic weight. Let it be remarked at once that to-day, by more refined observations on light and X-rays, we can obtain far fuller enlightenment on atomic structure. Thus we have in the first place the phenomenon of dispersion of the light passing through, which gives information on the binding forces of the electrons; or again, the interference rings which careful observations reveal in the scattered radiation from X-rays

(p. 68), and which supply data on the diameter of the electronic envelope of the atom, and on the distance between the centres of the atoms in molecules (Debye, 1929).

We pass now to the attempts to investigate atomic structure by means of scattering experiments on beams of electrons. These were first made by Lenard and his collaborators, who used cathode rays. We have seen in Chapter I that the diameter of an atom is

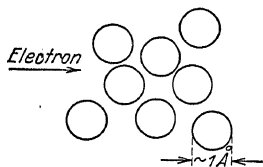


Fig. 2.—Passage of electrons through matter; according to the conceptions of the kinetic theory of gases the molecules are structures of size  $\sim 1 \text{ \AA}$ ; if they were impenetrable to an electron, the latter could not push its way through thin foil.

of the order of magnitude  $10^{-8} \text{ cm}$ . If the atoms were massive spheres, as in the scheme of fig. 2, then collisions with these spheres would necessarily very soon stop a cathode ray particle completely. The systematic investigations of Lenard, however, gave a precisely contrary result. It was found, in fact, that atoms are almost perfectly transparent to swift electrons. This suggested to Lenard (1903) the idea that an atom consists of a very small impenetrable centre, which he called a *dynamid*; this centre is surrounded by an electron cloud of loose texture, offering scarcely any resistance to incident cathode rays. As

author of this hypothesis, Lenard may properly be credited with the first suggestion of the modern model of the atom.

Usually, however, the title of father of the atomic theory is given to Ruther-

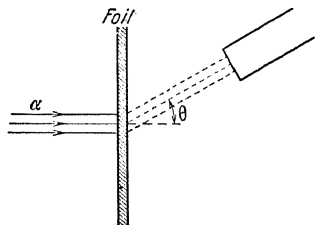


Fig. 3.—Arrangement for observing the scattering of  $\alpha$ -particles; a Geiger counter is used to count the number of  $\alpha$ -particles deflected from the primary direction through a definite angle  $\theta$ .

ford, who took up the research with more adequate instrumental resources, and carried it farther; to him we owe our concrete, quantitative ideas on atomic structure. For his experiments on scattering Rutherford (1911) used, not, like Lenard, the relatively light electron, but the much more massive  $\alpha$ -particle (fig. 3). On account of their greater mass,  $\alpha$ -particles are not noticeably deflected by the electrons in the atom, and therefore record only collisions with the more massive particles. The comparative penetrating effect of electrons and  $\alpha$ -particles may be illustrated roughly by the difference between a light rifle bullet and a heavy shell.

In the first place, Rutherford's experiments show definitely (1913)

that an atom, except for a small massive nucleus, is almost perfectly empty—a result which had already been found by Lenard. Since the  $\alpha$ -particles are perceptibly deflected by the nucleus only, we can deduce the law of deflection from the law of distribution of the  $\alpha$ -particles which have been scattered by a piece of foil. The definite result was found that the effective deflecting force is the Coulomb force  $2Ze^2/r^2$ , where  $2e$ , as we

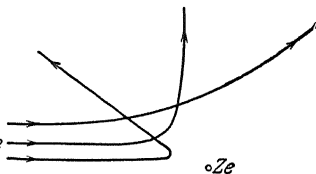


Fig. 4.—Scattering of  $\alpha$ -particles by a nucleus of charge  $Z$ ; the paths of the  $\alpha$ -particles are hyperbolas.

know, is the charge of the  $\alpha$ -particle, and  $Ze$  is the charge of the nucleus. The paths of the  $\alpha$ -particles are hyperbolas with the nucleus as focus (fig. 4).

If the incident ray contains one  $\alpha$ -particle per square centimetre, then according to Rutherford (Appendix IX, p. 276) the number of particles per unit solid angle which suffer a deflection  $\phi$  is given by

$$w(\phi) = Z^2 \left( \frac{m}{M} \right)^2 \left( \frac{e^2}{mc^2} \right)^2 \frac{1}{(v/c)^4 \sin^4 \frac{1}{2} \phi}$$

( $m$  = mass of electron,  $M = 4m_{\text{H}}$  = mass of  $\alpha$ -particle); here again the effective cross-section of the electron occurs as a factor, but the electronic radius  $e^2/mc^2$  is multiplied by the small factor (characteristic of the dimensions of the nucleus)

$$\frac{m}{M} = \frac{m}{4m_{\text{H}}} = \frac{1}{4 \times 1840} = 1.36 \times 10^{-4}.$$

The experiments showed that the number  $Z$ , which gives the nuclear charge, is equal to the number which would be assigned to the element in question in a consecutive enumeration of the elements in the periodic system. If the atom is to be neutral, the number giving the nuclear charge must agree with the number of electrons in the electron cloud surrounding the nucleus, as determined by optical and X-ray scattering experiments. The chemical behaviour of the atom depends of course on the external electrons, so that it is not the mass of the atom, but its atomic number (or number giving the nuclear charge) which determines its chemical properties. Isotopes have the same atomic number.

For collisions which are nearly central, i.e. for scattering through wide angles, deviations occur from the distribution of the scattered  $\alpha$ -particles determined by Coulomb's law. From this we must infer

that Coulomb's law only holds down to distances of about  $10^{-13}$  cm. The nucleus also has a finite size; it is worthy of remark that the "nuclear radius" is of the same order of magnitude as the radius of the electron.

According to Rutherford, the nuclear atom may be described as follows. In the centre of the atom there is the nucleus; this was thought of as composed of  $n_+$  protons and  $n_-$  electrons, its charge being therefore equal to

$$Ze = (n_+ - n_-)e,$$

and its mass, apart from the mass defect to be discussed later (p. 55), equal to  $n_+$  referred to  $H = 1$ . Possible isotopes must have the same nuclear charge number  $Z$ , and can therefore only consist of

$$n_+ \pm x \text{ protons}$$

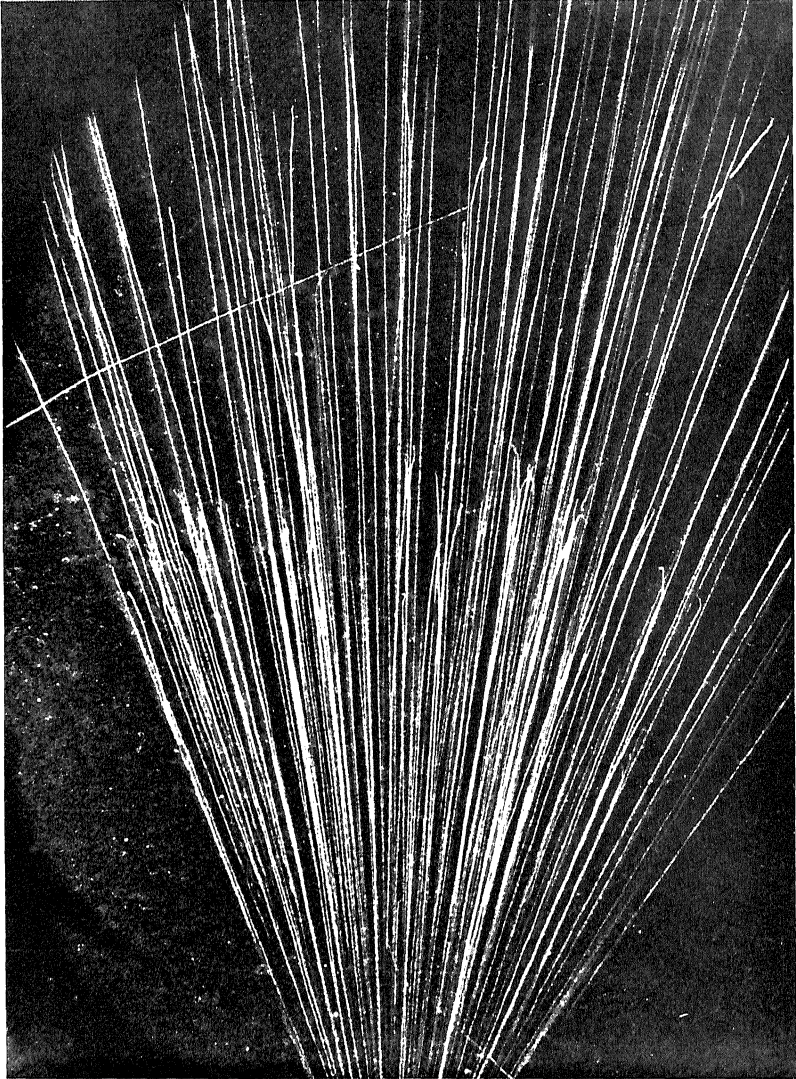
and

$$n_- \pm x \text{ electrons.}$$

Since the discovery of the neutron and the positron, there are clearly other possible ways in which the nucleus may be thought of as built up. Only one of these models, however, has been found feasible, viz. a nucleus composed of  $p$  protons and  $n$  neutrons; the nuclear charge number (atomic number) is then  $Z = p$ , the atomic mass number is  $A = p + n$ . The considerations in favour of this model will be presented immediately (§ 4, p. 57).

Round the nucleus, as has been mentioned, there move in the neutral atom  $Z$  electrons, which fill a sphere of radius  $\sim 10^{-8}$  cm. To get an idea of the dimensions, and the emptiness, of an atom, take the following illustration. If we imagine a drop of water to be expanded to the size of the earth, and all the atoms in it also enlarged in the same proportion, an atom will have a diameter of a few metres. The diameter of the nucleus, however, will be only something like  $1/100$  mm., and this must also be the order of magnitude of the  $Z$  electrons which fly round in the sphere corresponding to the atom.

It follows from these relations of magnitude that in the great majority of physical and chemical processes the nuclei act simply as positively charged point masses; only the external electronic system is essential. Research has therefore been devoted to this first, with such success indeed that we are to-day in possession of a theory which seems to be in complete accord with experiment. With this we shall deal in detail later. Nuclear physics, on the contrary, is still in its initial stages. Experimental research, it is true, advances rapidly, but



Ch. III, Fig. 6. Disintegration of a nitrogen nucleus by an  $\alpha$ -particle. A proton of very long range is emitted.  $\alpha$ -particles of two different ranges are seen, emitted by a mixture of thorium B and C. (From *Proc. Roy. Soc. A*, Vol. 136.)



theory does not as yet keep pace with it. And yet the problem of investigating the structure of nuclei, and of reducing to simple principles the typical laws and relations thereby brought to light, is without doubt the central problem of all natural science. We must content ourselves here with giving an account of the experimental methods and results. Our present theoretical resources are capable of dealing with only a few features of the case (§ 9, p. 176).

#### 4. Mass Defect and Nuclear Binding Energy.

We have already mentioned that the mass of a nucleus is not an exact multiple of the mass of the proton. This is really a matter of course, for, in order to remove a proton, neutron or  $\alpha$ -particle from the nucleus of a higher element, we must supply energy (except in radioactive substances, in which such disintegrations can occur spontaneously, i.e. without previous supply of energy). This loss of energy when elementary particles become united to a nucleus is, however, according to Einstein (1905), equivalent to a loss of mass, so that the final product is lighter than the sum of the weights of the individual components in the separated state. Take a numerical example. The greatest deficit occurs in the formation of a helium atom from four hydrogen atoms. The latter have the total mass  $4 \times 1.008$ , and in the combined state (helium) the mass 4.004; the energy freed in the process of combination is therefore, by Einstein's formula,

$$\begin{aligned} E &= (4 \times 1.008 - 4.004)c^2 = 0.028 c^2 \\ &= 0.25 \times 10^{20} \text{ ergs/mole.} \end{aligned}$$

In order to split up a helium atom into four hydrogen atoms, at least this amount of energy must be supplied. To make it easier to grasp its order of magnitude, convert ergs into calories; we find the value of  $E$  to be the enormous one of  $6.4 \times 10^8$  kilocalories per mole. For the sake of comparison, it may be noted that heats of combustion are of the order of magnitude of a few hundred kilocalories per mole. In nuclear physics it is customary to quote mass defects or binding energies simply in atomic weight units, e.g. for He,  $E = 0.028$ ; or, since they are small fractions of the unit, in multiples of the electronic mass  $m$  (or the electronic energy  $mc^2$ ), so that for He,

$$.E = 0.028 \times 1840 mc^2 = 52 mc^2.$$

We shall frequently make use of the unit  $mc^2$ . It corresponds to a kinetic energy of the electron, requiring about 500,000 volts to produce; we say that it is equal to  $5 \times 10^5$  electron volts ( $eV$ ).

The measurements of mass defects carried out with very great accuracy by Aston (since 1920) supply information, therefore, on the " heats of formation " of nuclei, i.e. on the energy relations in nuclear construction. These are found to follow certain perfectly definite rules. It has proved to be convenient to divide up the various nuclei into four groups,

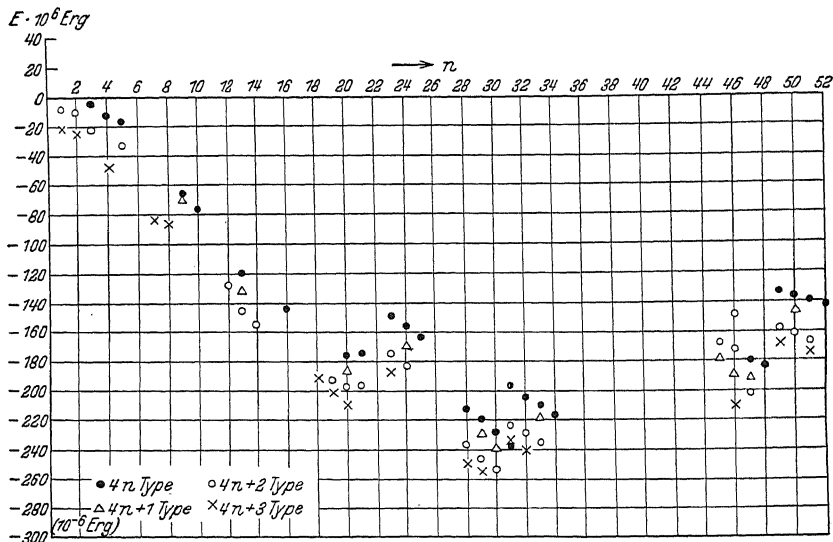


Fig. 5.—Binding energies (mass defects) in ergs, plotted against  $n$ , where  $n$  is the greatest integer such that  $4n$  does not exceed the atomic weight

(From Proc. Roy. Soc. A., Vol. 136)

according as their mass numbers are divisible by 4, or on division by 4 leave a remainder of 1, 2 or 3; we shall return to this point below (p. 58). The binding energies (and therefore the mass defects) are shown in fig. 5; it will be seen that as we ascend in the periodic system the binding energies (mass defects) increase rapidly at first, then slowly fall off again as we come nearer to the radioactive elements.

The difficulties of a theoretical explanation of these facts arise not only from our ignorance of the forces which act between particles at distances approaching values smaller than the electronic radius, but from certain other circumstances which cannot be understood without explanations to be given in later chapters. One point is evident at once. Since the diameter of the nucleus has been shown by collision experiments to be of the same order of magnitude as the diameter of the electron, it follows that the obviously suggested and long-accepted idea that nuclei are built up of protons and electrons alone,



presents difficulties. There is no room, as it were, for the electrons inside a nucleus; indeed there are many indications that the radius of the proton is approximately equal to that of the electron, and not 1840 times smaller, as we should expect from the formula  $e^2/mc^2$ . The conclusion that there are no free electrons in nuclei can be drawn more definitely from measurements of the angular momentum of nuclei, the so-called nuclear spin. We shall return to the matter again (§ 9, p. 176); here let only this much be said: in the quantum theory the angular momentum of a particle is always a multiple of an elementary unit; and if several particles are conjoined, there must be simple whole number relations between the angular momenta of the individual partners and of the whole. These relations, though in the domain of the external electronic envelope they hold without exception, are not fulfilled in nuclei—a specially well-verified case is the nitrogen nucleus  $N^{14}$ . If we construct the nucleus from protons and electrons, then the angular momenta behave as if there were only protons present (Heitler, Herzberg, 1929).

Now we have to-day other particles at our disposal, above all the neutron, and it is an obvious suggestion to try the experiment of calling upon these for the elucidation of nuclear structure. All the above difficulties are avoided, if we assume that the nuclei are composed of protons and neutrons (Iwanenko, Heisenberg, 1932). In the first place, with the help of simple assumptions about the interaction forces we can then understand the fundamental fact that in light atoms the atomic weight  $A$  is double the atomic number  $Z$ , while in heavy atoms it increases somewhat faster than  $2Z$ . Between two protons there acts the Coulomb repulsion of their positive charges; between two neutrons there presumably acts only a very small force, which may be neglected. Between a neutron and a proton there must be a very considerable attraction, to make the formation of the nucleus possible. We can picture this force in some such way as we picture chemical binding forces, which act between atoms, and whose quantum-theoretical mechanism we know. It rests upon exchange of electrons (p. 255; Appendix XXXI, p. 340) between two atoms; one or two electrons oscillate from one to the other, and so produce a binding between a pair of atoms, which has the character of saturation (chemical valency), while other atoms are only attracted very much more weakly. In the same way we can imagine an exchange of charge between proton and neutron, connected with a kind of valency attraction. The numbers  $p$  of the protons, and  $n$  of the neutrons, so adjust themselves that as much energy as possible is set free in the binding process. This means,

if we neglect the Coulomb repulsion between the protons, that as many pairs proton-neutron are formed as possible. If now the atomic weight, or more exactly the mass number,  $A = p + n$  is given, the most stable state for small numbers is that in which  $p = n$ ; since  $p = Z$ , the atomic number, it follows that  $A = 2Z$ . The more protons, however, there are in the combination, the more will the Coulomb repulsion come into account, and that clearly to the disadvantage of the protons; hence the difference  $n - p$  becomes  $> 0$ , and increases with increasing  $p = Z$ . It follows that  $A = n + p > 2p$ , or  $A > 2Z$ , and that  $A - 2Z$  increases as  $Z$  increases; and this is actually the case. It is very probable that two pairs, proton-neutron, within a nucleus always combine so as to form an  $\alpha$ -particle, so far as that is possible; for we know this binding to be particularly rigid (the mass defect is great; see p. 55); also, it is  $\alpha$ -particles which are emitted from the unstable (radioactive) nuclei; and, finally, the nuclei with atomic weight divisible by 4 form a specially regular series of mass defects (p. 56). Moreover, the stability of the  $\alpha$ -particle can be understood, viz. on the ground of the behaviour of the angular momenta or spins of the elementary particles, which have a tendency towards saturation (formation of systems with angular momentum nil, in accordance with Pauli's principle; see § 5, p. 159). This theory gives a very satisfactory account of radioactive  $\alpha$ -disintegration. As for  $\beta$ -disintegration, the expulsion of an electron, the new conception has the advantage of reducing this to a single elementary process, viz. the disintegration of the neutron into a proton and an electron. Recently, the expulsion of positrons has been observed from short-lived (artificially produced) nuclei (p. 64); this process would be the disintegration of the proton into neutron + positron. The characteristic of the ordinary  $\beta^-$ -emission is, as we have seen (p. 31), a continuous velocity spectrum; the same is true of the  $\beta^+$ -emission also. We have already pointed out the fundamental difficulty arising from this fact; it means that the principle of energy is infringed in these processes (as also, for that matter, the principle of the conservation of angular momentum, or spin). To save it, there is nothing for it but to assume a third particle, which is concerned in the process, but which in consequence of its properties escapes direct observation (Pauli, Fermi, 1934). It must be very light, at least as light as the electron, and can have no charge; it is therefore called a "neutrino". This is no mere playful suggestion; for although the traces of the neutrino will perhaps never be directly demonstrable, yet there is an indirect proof of its existence, quite analogous to the one for the neutron

(§ 6, p. 40). We see the tracks of atomic nuclei which have been struck; we then infer the existence of the neutron by applying the principle of conservation of momentum. In the case of the neutrino such hits will be too rare; in their stead, however, we can use the recoil at the emission. In ordinary radioactive materials, whose atoms are all very heavy, the track of the recoiling atomic residue at a  $\beta$ -emission is certainly too minute; but to-day we can also manufacture light radioactive atoms artificially (§ 6, p. 61), and with these the recoil would certainly be observable. If then the direction of the tracks of the electron and of the atomic residue are not exactly opposite, the assumption of a third invisible particle would be preferable to giving up even the conservation of momentum (Bethe, Peierls, 1934). If, besides, experiments were to lead to a definite unique mass (as calculated from energy and momentum, see Appendix V, p. 271), then the existence of the neutrino could not be doubted. These experiments, however, have not yet been carried out.

### 5. Heavy Hydrogen and Heavy Water.

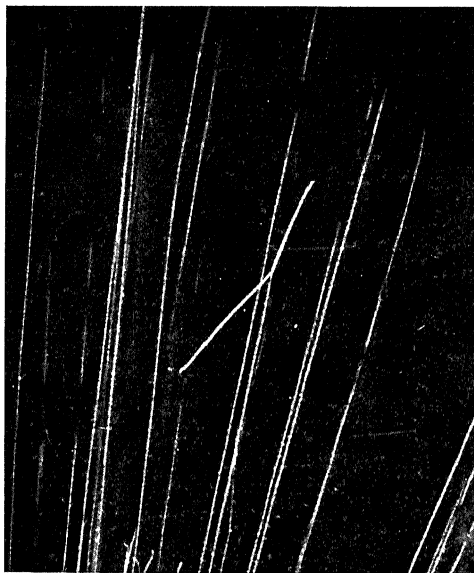
Most isotopes play no part in the everyday practice of physicist and chemist, since their differences in mass are relatively slight, and their properties almost identical. It is different with the recently discovered isotope of hydrogen, the nucleus of which has about double the mass of the proton; it has been given the name *deuteron* (the second) and the symbol D. In this case essential differences occur in physical and chemical behaviour, so that we can speak of an actually new element, which is called *deuterium*. The importance of this discovery becomes particularly evident when we remember that the most important compound of hydrogen, water ( $\text{H}_2\text{O}$ ), which physics from the beginning has taken as a standard substance, now turns out to be a mixture of several kinds of molecule ( $\text{H}_2\text{O}$ ,  $\text{HDO}$ ,  $\text{D}_2\text{O}$ ), the molecular weights of which differ from each other by 5 and 10 per cent. It is therefore fortunate that, for other reasons of a practical nature, we have long since given up the definition of the kilogram as the mass of 1000 c.c. of water at its maximum density, defining it instead as the mass of a certain piece of platinum-iridium. We may relate the almost romantic story of the discovery, which reminds one of the discovery of argon in the atmosphere by Rayleigh and Ramsay (1894); in both cases the key to the discovery was belief in the reality of minute discrepancies in different measurements of the same magnitude. To begin with, two isotopes of oxygen, of masses 17 and 18, were discovered spectroscopically; since

the frequencies of vibrations are inversely proportional to the square root of the mass of the vibrating particle, the isotopes betrayed their presence by the occurrence of displaced lines in the molecular spectrum. However, the intensity of the lines was very weak, and the amounts of the isotopes so minute that they could not affect practical measurements. Theoretically, however, the demonstration of their existence, and the measurement of their relative amounts, had important consequences. From these amounts, and the chemical molecular weights, we can in fact calculate the mass of the H-atom, referred to the principal isotope of O as 16. On the other hand, this same magnitude has been determined by Aston with the mass-spectrograph. It was found by Birge and Menzel (1931) that a difference of 1/5000 remained over, and they concluded that hydrogen also must contain a small amount, about 1/4000, of a heavy isotope, of mass 2.\* Thereupon spectroscopic investigations on hydrogen were undertaken by Urey, Brickwedde and Murphy (1931). We shall see that the position of the atomic lines of an element also depends on the mass of the nucleus (p. 97). Actually, weak satellites of the principal hydrogen lines were found at the right distance. Various methods were then applied to produce enrichment in the heavy isotope D, the most successful of these being ordinary electrolysis (Urey and Washburn, Lewis and Macdonald, 1932). It was found that the light  $H_2$  escapes from the cathode 5 or 6 times faster than the heavy  $D_2$ . If the original concentration was 1 in 5000, the electrolysis of 6 litres of ordinary water should give us about 1 c.c. of pure heavy water. In point of fact, sufficient quantities of heavy water are now available to allow its properties to be investigated in all directions. Another method of separation, that of Hertz already mentioned (p. 37), depends on the difference in the rates of diffusion of the gases; by this method about 1 c.c. of  $D_2$  was obtained of such purity that in the spectrum the atomic lines of H (the so-called Balmer lines; see p. 93) were no longer visible at all. The exact mass of the D-atom is 2.0147 (Aston, 1936), referred to the principal isotope of O as 16, while that of the H-atom is 1.0081. From the point of view taken by Heisenberg (p. 57), we must look upon the D-nucleus as the combination proton + neutron. This has been directly proved by decomposition into these particles by means of  $\gamma$ -rays (Chadwick and Goldhaber, 1934; p. 65). The properties of heavy water  $D_2O$ , differ very noticeably from those of ordinary water,  $H_2O$ ; the freezing-point is about  $3.8^\circ$ , the boiling-point about  $1.4^\circ$ , higher; the

\* Recent measurements of Aston show, however, that this conclusion which has proved so fruitful was based on incorrect measurements.



Ch. III, Fig. 7.—Elastic collision of an  $\alpha$ -particle with a nucleus of hydrogen. The tracks of the two particles after the collision are of nearly equal length.



Ch. III, Fig. 8. Elastic collision of an  $\alpha$ -particle with a nucleus of nitrogen. The short track corresponds to the nitrogen. (From *Proc. Roy. Soc. A*, Vol. 134.)



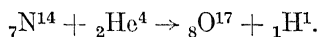
density is as much as about 11 per cent greater than that of  $\text{H}_2\text{O}$ . Similar results hold for other combinations in which H is wholly or partly replaced by D. The velocities of reaction in these show considerable changes. Thus an entirely new branch of chemistry has emerged, which is also important for biology.

For the physicist, however, the importance of heavy hydrogen is mainly due to the fact that we can produce D canal rays just as easily as H canal rays, and accordingly have now a new means of bombarding other nuclei, which has already given us valuable information about their constitution. Of this we shall speak in the following section.

### 6. Artificial Disintegration of the Nucleus and Artificial Radioactivity.

As we know, all experience shows that radioactive nuclear disintegration cannot be affected by any ordinary physical means; it takes place spontaneously in accordance with statistical laws. If we had to depend on these processes alone, we would never get direct information as to the structure of the majority of (non-radioactive) nuclei, and would always have to rely on hypotheses.

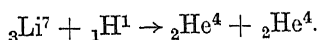
The greatest importance, therefore, attaches to Rutherford's discovery (1919) that nuclei can be broken up by bombardment with  $\alpha$ -rays. The first element with which this was accomplished was nitrogen; in the cloud chamber, nuclear hits are occasionally seen in which the track of the incident  $\alpha$ -particle suddenly disappears, and the track of the nucleus which was struck, and that of a particle of greater range, begin where the first track ends (figs. 6, 7, 8, Plates III, IV). The same result was afterwards obtained with many other elements, boron, fluorine, neon, sodium, &c. (Rutherford, Chadwick, Kirsch and Petterson, 1920-1925), and it was shown by deflection experiments that the light particles expelled are protons. The  $\alpha$ -particle is obviously captured by the nucleus; from  ${}_7\text{N}^{14}$  (below, the atomic number  $Z$ ; above, the mass number or rounded off atomic weight  $A$ ) is formed the isotope of oxygen of mass 17, already mentioned above, in accordance with the equation



We are here at the starting-point of a sort of nuclear chemistry, the laws of which, as in the example just given, can be written down in a form exactly analogous to chemical formulæ. (We can also add the binding energies, which are analogous to heat evolved or absorbed, and which must exactly compensate the kinetic energies of the particles.)

The phenomenon, contrary to what was thought at first and is still frequently said, is not one of destruction of the nucleus, but of a nuclear transformation, which should rather be called nuclear construction. More recently, many other swift corpuscular rays have been used to bombard the nucleus, and the resulting transformations studied.

The first step was the discovery by Cockcroft and Walton (1932) that artificially produced beams of protons, with energy of 120,000 electron volts, equal to about  $\frac{1}{2}mc^2$ , a value quite trifling compared with the energy of natural  $\alpha$ -rays (about  $16 mc^2$ ), are capable of breaking up the nucleus of the lithium atom, according to the relation



In point of fact, it was shown in the Wilson chamber that two helium nuclei, i.e.  $\alpha$ -particles, were always shot out simultaneously in opposite directions from the bombarded lithium foil (Kirchner, Dee and Walton, 1934; see Plate V, fig. 9). Similar transformations were successfully brought about in a whole series of nuclei, and it is remarkable that the measurements of the masses of the nuclei involved with help of the mass-spectrograph and the determination of the kinetic energies of the particles before and after the disintegration can be performed with such accuracy that one gets a direct experimental verification of Einstein's law  $E = mc^2$  (Bainbridge, 1933).

The nuclear-chemical processes take place without the incident proton approaching the nucleus very closely in virtue of its kinetic energy; for of course we can calculate how near the nuclei come by equating the potential energy  $e^2Z/r$  at the minimum distance to the kinetic energy of a proton  $eV$ , where  $V$  is the accelerating potential: thus

$$V = \frac{eZ}{r}, \quad r = \frac{eZ}{V}.$$

Hence, with an energy  $eV$  of  $10^5$  electron volts, i.e.  $1.59 \times 10^{-7}$  ergs, so that  $V = \frac{1}{3} \times 10^3$  e.s.u., we have, for  $Z = 1$ ,

$$r = 4.77 \times 10^{-10} \times 1 \times 3 \times 10^{-3} = 14 \times 10^{-13} \text{ cm.},$$

which is of the order of magnitude of the nuclear radius (§ 3, p. 54). But the disintegration effect in lithium is demonstrable down to less than 30,000 electron volts (Rausch von Traubenbergs and Döpel, 1933), in which case the proton is still far outside the nucleus. The explanation of this fact will be given later (p. 183).

The heavy hydrogen isotope D has also been used to bring about disintegration (Lewis, Livingstone and Lawrence, 1933) and proved

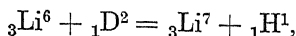


TABLE III  
 ARTIFICIAL NUCLEAR TRANSFORMATIONS

 Radioactive nuclei emitting  $\begin{cases} \text{electrons} \\ \text{positrons} \end{cases}$  are indicated by  $\begin{cases} - \\ + \end{cases}$ .

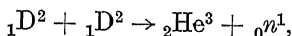
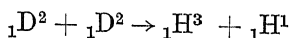
Bombarding Particle Bombarded Nucleus	$\gamma$	${}^0n^1$	${}_1H^1$	${}_1D^2$	${}_2He^4$
${}_1D^2$	${}_1H^1 + {}^0n^1$			${}_1H^3 + {}_1H^1$ ${}_2He^3 + {}^0n^1$	
${}_3Li^6$		${}_2He^4 + {}_1H^3$	${}_2He^4 + {}_2He^3$	${}_3Li^7 + {}_1H^1$ ${}_2He^4 + {}_2He^4$	
${}_3Li^7$		${}_3Li^8 + \gamma (?)$	${}_2He^4 + {}_2He^4$	${}_2He^4 + {}^0n^1$ ${}_3Li^8 + {}_1H^1$	${}_4Be^{10} + {}_1H^1 (?)$
${}_4Be^9$	${}_4Be^8 + {}^0n^1$ (?)	${}_2He^6 + {}_2He^4$ (?)	${}_4Be^8 + {}_1D^2$	${}_5B^{10} + {}^0n^1$	${}_6C^{12} + {}^0n^1$
${}_5B^{10}$		${}_3Li^7 + {}_2He^4$ ${}_2He^4 + {}_1H^3$		${}_6C^{11} + {}^0n^1$ ${}_3He^4$ ${}_5B^{11} + {}_1H^1$	${}_7N^{13} + {}^0n^1$ ${}_6C^{13} + {}_1H^1$
${}_5B^{11}$			${}_3He^4$	${}_6C^{12} + {}^0n^1$ ${}_5B^{12} + {}_1H^1$ ${}_3He^4 + {}^0n^1$	
${}_6C^{12}$		${}_3He^4 + {}^0n^1$ ${}_4Be^9 + {}_2He^4$	${}_7N^{13} + \gamma$	${}_7N^{13} + {}^0n^1$	
${}_6C^{13}$				${}_7N^{14} + {}^0n^1$ ${}_5B^{11} + {}_2He^4$	
${}_7N^{14}$		${}_5B^{11} + {}_2He^4$ ${}_6C^{14} + {}_1H^1$ ${}_3Li^7 + {}_2He^4$		${}_7N^{15} + {}_1H^1$ ${}_6C^{12} + {}_2He^4$	${}_8O^{17} + {}_1H^1$ ${}_9F^{17} + {}^0n^1$
${}_8O^{16}$		${}_6C^{13} + {}_2He^4$			
${}_9F^{19}$		${}_7N^{16} + {}_2He^4$	${}_8O^{16} + {}_2He^4$	${}_9F^{20} + {}_1H^1$	${}_{10}Ne^{22} + {}_1H^1$ ${}_{11}Na^{22} + {}^0n^1$
${}_{10}Ne^{20}$					${}_{11}Na^{23} + {}_1H^1$
${}_{11}Na^{23}$		${}_{11}Na^{24} + \gamma$		${}_{11}Na^{24} + {}_1H^1$	${}_{12}Mg^{26} + {}_1H^1$ ${}_{13}Al^{26} + {}^0n^1$
${}_{12}Mg^{24}$		${}_{11}Na^{24} + {}_1H^1$			${}_{13}Al^{27} + {}_1H^1$ ${}_{14}Si^{27} + {}^0n^1$
${}_{13}Al^{27}$		${}_{11}Na^{24} + {}_2He^4$ ${}_{12}Mg^{27} + {}_1H^1$		${}_{13}Al^{28} + {}_1H^1$	${}_{14}Si^{30} + {}_1H^1$ ${}_{15}P^{30} + {}^0n^1$

extremely effective. Thus, e.g., we get (according to Cockcroft and Walton, 1934) a transformation of the light lithium isotope into the heavy one,

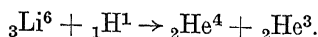


and the protons so obtained have a very long range. There is no need to quote all the transformations so far effected; with other processes, to be discussed immediately, they are collected in Table III (p. 63).

An important discovery is that of new isotopes of hydrogen and helium, viz.  ${}_1\text{H}^3$  (which should be called triton) and  ${}_2\text{He}^3$ . They were first found in disintegration experiments in the Cavendish Laboratory, corresponding to the reactions (Oliphant, Harteck, Rutherford, 1934)

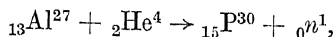


and (Oliphant, Kinsey and Rutherford, 1933)

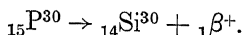


Soon afterwards  ${}_1\text{H}^3$  was directly demonstrated by the mass-spectrograph as a weak admixture with  ${}_1\text{H}^1$  and  ${}_1\text{D}^2$  (Bleakney, Lozier and Smith, 1934).

Neutrons were next employed as bombarding particles (Feather, Harkins, 1932). It is probable on the face of it that these should give a large output, since they are uncharged and so not repelled by the field of the struck nucleus. This surmise has been confirmed, and led to the discovery of a new class of transformations, of which we shall speak presently. In point of fact, these transformations, which may be described as an artificial production of radioactivity, were discovered not with neutrons, but with  $\alpha$ -rays, the discoverers being Irène Curie and her husband Joliot (1934). For instance, it was found that aluminium which was bombarded with  $\alpha$ -rays and so sent out neutrons continued to radiate after the irradiation was stopped, the emission consisting of positive electrons. Clearly the initial product must, in accordance with the equation



be an isotope of phosphorus, which is radioactive, and passes over into an isotope of silicon, with emission of a positron:



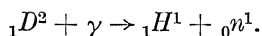
The half-value period of the "radio-phosphorus" was found to be

3 min. This is long enough to allow the correctness of the interpretation to be tested in these and similar processes by the old method of chemical precipitation (p. 36). The proton and its isotope D also proved useful for the production of radioactive nuclei.

The most abundant output, however, comes from neutrons, with which Fermi and his collaborators (1934) have produced new radioactive nuclei from the majority of known atoms. In spite of the rather short lifetimes of these nuclei (some minutes) it was possible to determine the chemical character of the products, and it was found that the atomic number of the nucleus which captures a neutron changes by 2 or 1 or 0, corresponding to the emission of an  $\alpha$ -particle, a proton or of light.

By passing them through substances like water or paraffin wax which contain many hydrogen atoms, neutrons are slowed down by the impacts with the protons conferring on these particles of equal mass half their energy at each collision. In this way it is possible to have neutrons nearly in thermal equilibrium with the surrounding substance. These slow neutrons are particularly effective in producing new nuclei, and they show many interesting properties, such as selective absorption in special elements. The study of these processes is one of the most powerful methods for revealing the laws of nuclear structure.

Recently Chadwick and Goldhaber (1934) discovered a new type of disintegration, using  $\gamma$ -rays as bombarding agent. The heavy hydrogen isotope  $D$  was shown to break up according to the relation



Thus the binding energy of the  $D$ -nucleus ( $D = p + n$ ) was determined to be about  $4.5mc^2$ , and a value for the mass of the neutron could be derived (1.0090).

At this point we conclude our brief survey of the state of knowledge with regard to nuclei, and now turn to the structure of the external electronic system. We are thus brought into regions whose linear dimensions are several thousand times greater than those of the nucleus. Here, by the co-operation of experimental and theoretical research, the fundamental laws have in all essential respects been made clear. But this was only made possible by the renunciation of ideas to which people's minds had become so accustomed from their experience of events on the large scale that a profound critical analysis had first to be carried out.

## CHAPTER IV

### Wave-Corpuscles

#### 1. Wave Theory of Light. Interference and Diffraction.

The ideas which we have arrived at in the preceding chapters with regard to the structure of matter all rest on the possibility of demonstrating the existence of fast-moving particles by direct experiment, and indeed of making their tracks immediately visible, as in the Wilson cloud chamber. These experiments put it beyond doubt that matter is composed of corpuscles. We are now to learn of experiments which just as definitely seem to be only reconcilable with the idea that a molecular or electronic beam is a *wave train*. Before we enter upon this, however, we shall briefly recall the main facts of wave motion in general, using the phenomena of optical diffraction as a concrete example.

While in the eighteenth century physicists almost universally adhered to Newton's emission theory (about 1680), according to which light consists of an aggregate of very small corpuscles, which are sent out by the source of light, and the wave theory of Huygens (1690) could claim only a few supporters (among them the great mathematician Euler), the state of matters changed completely when at the beginning of the nineteenth century Young made the discovery that in certain circumstances two beams of light can enfeeble each other, a phenomenon quite incapable of explanation on the corpuscular theory. The results of the further investigations of Young and Fresnel spoke unequivocally in favour of the wave conception of Huygens, for it is impossible to explain interference phenomena except by a wave theory.

We give here a short discussion of Young's *interference experiment* (fig. 1). The source of monochromatic light Q illuminates the double slit in the diaphragm B with parallel light by means of the lens L. On the screen S behind the diaphragm a system of equidistant bright and dark strips (fringes) appears. How this comes about may be explained as follows. From the two openings in the diaphragm spherical waves spread outwards; these are "coherent", i.e. they are capable

of mutual interference. The two wave motions become superimposed, and reinforce each other at those places where a crest of the one wave coincides with a crest of the other; on the contrary, they destroy each other where a crest of the one wave is superimposed on a hollow of the other. Hence we can tell at once at what places on the screen there will

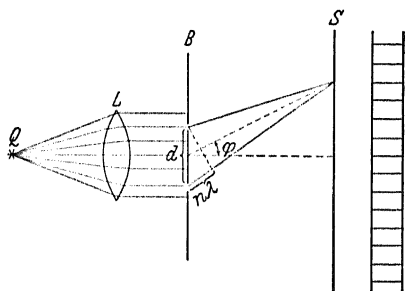


Fig. 1.—Diffraction at two narrow slits close to each other. The diffraction pattern consists of a system of equidistant bright and dark bands (fringes).

be brightness; they are the points whose distances from the two openings in the diaphragm differ exactly by an integral multiple of the wave-length. From fig. 1 we see that the difference of the distances is  $d \sin \phi$ ; there is therefore on the screen

$$\left. \begin{array}{l} \text{brightness, if } d \sin \phi = n\lambda \\ \text{darkness, if } d \sin \phi = (n + \frac{1}{2})\lambda \end{array} \right\} (n = 0, \pm 1, \pm 2, \dots),$$

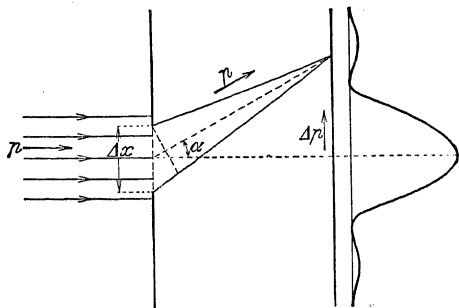
where  $d$  is the distance between the two openings, and  $\phi$  is the angle of deflection.

A similar diffraction pattern is also obtained when light passes through *one slit*. We can picture it roughly as due to the mutual interference of the elementary Huygens waves spreading out from the individual points of the slit. There are two essential differences, however, as compared with the previous case. In the first place, we easily see that the relation

$$d \sin \phi = n\lambda \quad (n = \pm 1, \pm 2, \dots),$$

where  $d$  is the slit-width, does not now give the places at which there is brightness, but those where there is darkness. For in the packet of wave trains which spread out from the slit in the direction given by the equation, all "phases" are in this case represented exactly the same number of times; i.e. we find in the packet exactly as many wave trains which reach the screen with a crest, as trains which arrive at it with a hollow; the trains will therefore extinguish each other. We find, further, that the diffraction maxima are not, as before, all equally bright; but that their intensity falls off very strongly from the middle maximum outwards, in the way indicated in fig. 2 by the wavy line shown at the side. It should also be specially emphasized

that *when the slit-width is reduced the diffraction pattern widens out*, as may easily be deduced either from the above equation defining the

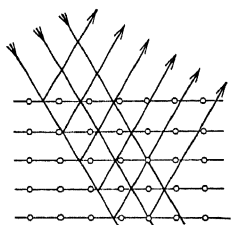


position of the diffraction minima, or directly from fig. 2.

The fact that the form of the diffraction pattern depends essentially on the

Fig. 2.—Diffraction at a slit. The diffraction pattern shows a strong maximum of intensity for the angle of diffraction  $\phi = 0$ , and also a series of equidistant maxima which become progressively weaker as the angle of diffraction increases.

wave-length of the light makes it possible to carry out spectral investigations by means of interference phenomena (ruled grating, echelon grating, Perot-Fabry plate, Lummer plate). For diffraction patterns to show themselves, it is necessary that the width of the slit employed should be of the order of magnitude of the wave-length of the light. If then we wish to obtain interference phenomena with X-rays, we have to use a grating in which the distance between the rulings is of the order of magnitude of  $1 \text{ \AA.} = 10^{-8} \text{ cm.}$



Such gratings are put into our hands by nature, as von Laue (1912) has shown, in the shape of crystals, in which the lattice distances are just of this order of magnitude. If a beam of X-rays is passed through a crystal, we do in fact obtain

Fig. 3.—Diffraction of X-rays at a crystal. As explained by Bragg, the rays are reflected at the lattice planes of the crystal, and thus made to interfere.

interference phenomena. Following Bragg (1913), we can interpret these as due to interference of the rays reflected at different lattice planes of the crystal (fig. 3). Moreover, Compton (1925) and others succeeded in producing X-ray interference in artificial gratings also, this being found possible at grazing incidence of the rays.

Interference of X-rays supplies a powerful weapon for investigating the structure of crystals. For this purpose we do not even require large pieces of the crystal, but can use it in the form of powder (Debye-Scherrer, 1915; Hull, 1917). The interference figures in the latter case are rings round the direction of the incident beam. Indeed, the powder grains may be of molecular size even. What is more, it is found that

the interference phenomena due to the individual atoms of the molecule are by no means completely obliterated by the irregular setting and motion of the molecules in liquids and gases. Circular interference rings are observed, from the intensity distribution of which we can draw conclusions with regard to the distances of the atoms in the molecule (Debye, 1929). However, we shall not enter further into these methods of investigating how matter is built up from atoms.

We add here, in the form of a scale, a summary of the wave-lengths of the types of radiation at present known (fig. 4). The scale is

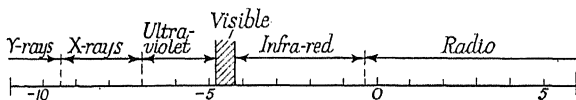


Fig. 4.—Logarithmic scale of wave-lengths; the numbers shown refer to a unit of length  $\lambda_0 = 1$  cm., and represent  $\log_{10}(\lambda/\lambda_0)$

logarithmic; the numbers shown are therefore indices of powers of 10. As unit  $\lambda_0 = 1$  cm. is taken. Next to the wide range of wave-lengths employed in wireless communication comes the region of infra-red waves, which affect our senses as radiant heat. Then comes the relatively narrow stretch of the visible region (7700 to 3900 Å.), followed by the ultra-violet and Schumann regions, which in turn lead into the domain of X-rays (100 to 0.05 Å.). The radioactive  $\gamma$ -radiation reaches to about 0.001 Å. The cosmic radiation seems to contain  $\gamma$ -rays of very short wave-length ( $10^{-5}$  Å.) which would lie beyond the scale of our figure.

## 2. Light Quanta.

Great as has been the success of classical ideas in the interpretation of interference phenomena, their incapacity to account for the processes of absorption and emission of radiation is no less striking. Here classical electrodynamics and classical mechanics absolutely fail.

To give a few examples of this failure, we recall the experimental fact that a hydrogen atom, for example, emits an infinite series of sharp spectral lines. Now the hydrogen atom possesses only a single electron, which revolves round the nucleus. By the rules of electrodynamics, an electron accelerated like this sends out radiation continuously, and so loses energy; in its orbit it would therefore necessarily get nearer and nearer the nucleus, into which it would finally plunge. The electron, which initially revolved with a definite frequency, will radiate light of this frequency; in the case when the

frequency of revolution changes (continuously) during the radiation process, it will also emit frequencies close to the ground frequency; but how then the spectrum of hydrogen should consist of a discrete series of sharply separated lines, it is quite impossible to understand.

Further, the stability of the atom is inexplicable. We may think by way of comparison of the system of planets circling round the sun, all, when undisturbed, moving in their fixed orbits. Suppose, however, that the whole solar system arrived in the neighbourhood of Sirius, for example; the mere propinquity would suffice to deflect all the planets out of their courses. If then the solar system moved away again to a distance from Sirius, the planets would now revolve round the sun in new orbits with new velocities and periods of revolution. If the electrons in the atom obeyed the same mechanical laws as the planets in the solar system, the necessary consequence of a collision of the atom with another atom would be that the ground frequencies of all the electrons would be completely changed, so that after the collision the atom would radiate light of wave-lengths also entirely different. In direct opposition to this, we have the experimental fact that an atom of a gas, which by the kinetic theory of gases is subjected to something like 100 million collisions per second, nevertheless, after these as before, sends out the same sharp spectral lines.

Finally, classical mechanics and statistics fail with the explanation of the laws of radiation of heat (or energy). We shall not go into this complex question in detail until later (Chap. VII, p. 185), and here merely quote the result to which Planck (1900) was led by the above considerations. To make the laws of radiation intelligible, he found the following hypothesis to be necessary: *emission and absorption of radiant energy by matter does not take place continuously, but in finite "quanta of energy"  $h\nu$  ( $h$  = Planck's constant  $6.55 \times 10^{-27}$  erg sec.,  $\nu$  = frequency).* On the other hand, the connexion with the electromagnetic theory of light is to be maintained to this extent, that the classical laws are to hold for the propagation of the radiation (diffraction, interference).

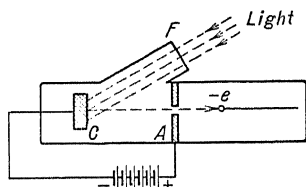
Einstein (1905), however, went even further than Planck. He not merely postulated quantum properties for the processes of absorption and emission of radiation, but also explained such properties as inherent in the nature of radiation itself. According to the *hypothesis of light quanta (photons)* which he advanced, light consists of quanta (corpuscles) of energy  $h\nu$ , which fly through space like a hail of shot, with the velocity of light. Daring as at first sight this hypothesis appears to be, there is nevertheless a whole series of experiments



which seem scarcely possible to explain on the wave theory, but which can be understood at once if we accept the hypothesis of the light quantum. Some account of these experiments, which were cited by Einstein himself in proof of his hypothesis, will now be given.

The most direct transformation of light into mechanical energy occurs in the photoelectric effect (Hertz (1887), Hallwachs, Elster and Geitel, Ladenburg). If short-wave (ultra-violet) light falls on a metal surface (alkalies) in a high vacuum, it is observed in the first place that the surface becomes positively charged (fig. 5); it is therefore giving off negative electricity, which issues from it in the form of electrons. We can now on the one hand, by capture of the electrons, measure the total current issuing

Fig. 5.—Production of photoelectrons (after Lenard). Light entering by the window F strikes the cathode C, and there liberates photoelectrons, which are accelerated (or retarded) in the field between C and A.



from the metal surface, and on the other hand determine the velocity of the electrons by deflection experiments or by a counter field. Exact experiments have shown that the velocity of the emergent electrons does not depend, as one might at first expect, on the intensity of the light; but that only their number increases as the light becomes stronger, the number being in fact proportional to the intensity of the light. The velocity of the photoelectrons depends only on the frequency  $\nu$  of the light; for the energy  $E$  of the electrons the following relation is found:

$$E = h\nu - A,$$

where  $A$  is a constant characteristic of the metal.

From the standpoint of the light quantum hypothesis, both these results can be understood at once. Every light quantum striking the metal and colliding with one of its electrons hands over its whole energy to the electron, and so knocks it out of the metal; before it emerges however, the electron loses a part of this energy equal in amount to the work,  $A$ , required to remove it from the metal. The number of electrons expelled is equal to the number of incident light quanta, and this is given by the intensity of the light falling on the metal.

Evidence even more patent than this, for the existence of light quanta, is given by the experiments of E. Meyer and W. Gerlach (1914) on the photoelectric effect with the small particles of metal dust; by

irradiation of these with ultra-violet light photoelectrons are again liberated, so that the metallic particles become positively charged. The advance on the previous case consists in this, that we can now observe the time relations in the process of charging the particles, by causing them to become suspended in an electric field, as in Millikan's droplet method for determining  $e$ , the elementary electric charge; a fresh emission of a photoelectron is then shown by the acceleration caused by the increase of charge.

If we start from the hypothesis that the incident light actually represents an electromagnetic alternating field, we can deduce from the size of the particles the time that must elapse before a particle of metal can have taken from this field by absorption the quantity of energy which is required for the release of an electron. These times are of the order of magnitude of some seconds; if the classical theory of light were correct, a photoelectron could in no case be emitted before the expiry of this time after starting the irradiation. But the experiment when carried out proved on the contrary that the emission of photoelectrons set in immediately the irradiation began—a result which is clearly unintelligible except on the basis of the idea that light consists of a hail of light quanta, which can knock out an electron the moment they strike a metal particle.

### 3. Quantum Theory of the Atom.

Planck's original quantum hypothesis was that to every spectral line there corresponds a harmonic oscillator of definite frequency  $\nu$  which cannot, as in the classical theory, absorb or emit an arbitrary quantity of energy, but only integral multiples of  $h\nu$ . Niels Bohr (1913) made the great advance of elucidating the connexion of these "oscillators" with one another and with the structure of the atom. He dropped the idea that the electrons actually behave like oscillators, i.e. that they are quasi-elastically bound. His leading thought was something like this. The atom does not behave like a classical mechanical system, which can absorb energy in portions which are arbitrarily small. From the fact of the existence of sharp emission and absorption lines on the one hand, and from Einstein's light quantum hypothesis on the other, it seems preferable to infer that the atom can exist only in definite *discrete stationary states*, with energies  $E_0, E_1, E_2, \dots$ . Thus only those spectral lines can be absorbed for which  $h\nu$  has exactly such a value that it can raise the atom from one stationary state to a higher one; the *absorption lines* are therefore defined by the equations  $E_1 - E_0 = h\nu_1, E_2 - E_0 = h\nu_2, \dots$ , where

$E_0$  is the energy of the lowest state in which the atom exists in the absence of special excitation.

If the atom is excited by any process, i.e. if it is raised to a state with energy  $E_n > E_0$ , it can give up this energy again in the form of radiation. It can in fact radiate all those quanta whose energy is equal to the difference of the energies of two stationary states. The emission lines are therefore given by the equation

$$E_n - E_m = h\nu_{nm}.$$

A direct confirmation of this theory can be seen in the following fact. If Bohr's hypothesis is correct, an excited atom has open to it various possible ways of falling back to the ground state by giving up energy as radiation. For example, an atom in the third excited state can either give up its excess of energy relative to the ground state in one elementary process by radiating a line of the frequency  $\nu_{30}$ ; or it can begin with a transition into the first excited state with the energy  $E_1$  and surrender of the energy quantum  $h\nu_{31}$ , and then in a second radiation process (frequency  $\nu_{10}$ ) fall back into the ground state; and so

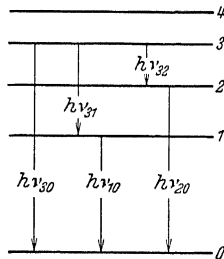


Fig. 6.—Ritz's Combination Principle. An atom in the third excited state can radiate its energy either in the form of a single light quantum of frequency  $\nu_{30}$ , or as two quanta, the sum of whose frequencies must be exactly  $\nu_{30}$ .

on (fig. 6). Since the total energy radiated must always be the same, viz.  $E_3 - E_0$ , the following relation must always exist between the radiated frequencies:

$$\nu_{30} = \nu_{31} + \nu_{10} = \nu_{32} + \nu_{20} = \nu_{32} + \nu_{21} + \nu_{10}.$$

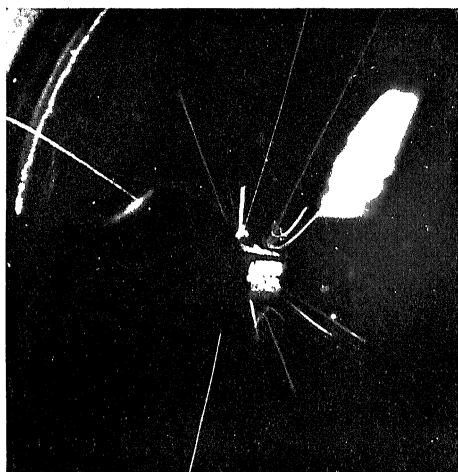
This *combination principle* must of course hold in all cases, and is a deduction from the theory which can easily be put to the test of experiment. Historically, it is true, the order of these two aspects of the matter was reversed; for Ritz, eight years before Bohr's theory was propounded, deduced this combination principle from collected spectroscopic material which had been obtained by experiment. It is by no means the case, however, that all possible "combination lines" do actually occur with perceptible intensity.

A further direct confirmation of Bohr's theory on the existence of discrete energy levels in the atom is given by the experiments of Franck and Hertz (1914). If the atoms are supplied with energy in any way, e.g. by electronic collision, i.e. by bombarding the atom with electrons,

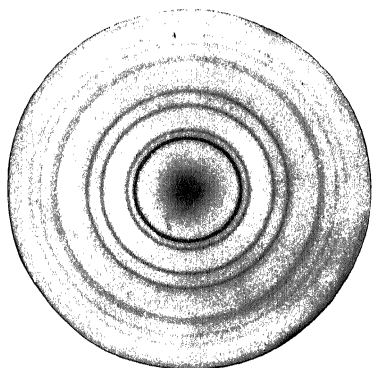
then the atoms can only take up such portions of energy as exactly correspond to an excitation energy of the atoms.

Thus, if we bombard the atoms with electrons whose kinetic energy is less than the first excitation energy of the atoms, no communication of energy from the colliding electron to the atom takes place at all (beyond the trifling amount of energy which is transferred in accordance with the laws of elastic collision, and shows itself only in the kinetic energy of the relative motion of the partners in the collision). With respect to collisions of no great strength the atoms are therefore stable in the ground state. Among such slight collisions are those which a gas atom is subjected to in consequence of the thermal motion of the particles of the gas. This is easily verified, roughly, as follows. The mean kinetic energy of a gas particle, by the results of the first chapter, is given by  $\bar{E} = \frac{3}{2}kT$ , where  $k = R/L = 1.37 \times 10^{-16}$  erg/degree; if the whole of this energy were converted at a collision into excitation energy, the energy quantum  $h\nu = \frac{3}{2}kT$  would be transferred to the particle struck; where, if we take  $T = 300^\circ$  K. (room temperature),  $\nu$  will be  $10^{13}$  sec $^{-1}$  approximately. On the other hand, frequencies of absorption lines in the visible or even in the infra-red part of the spectrum have values about  $10^{14}$  to  $10^{15}$  sec $^{-1}$ . Higher temperatures will therefore be necessary before "thermal excitation" of the atoms of the gas becomes possible.

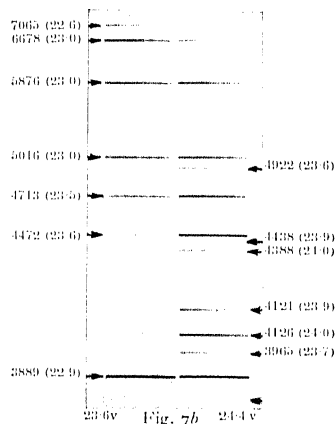
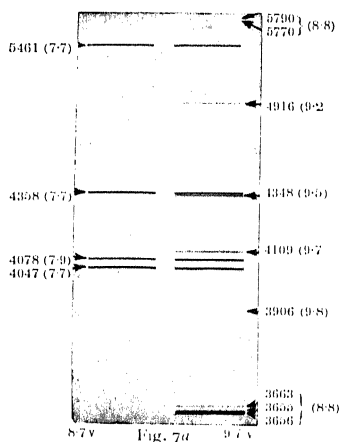
We return now to the collision experiments of Franck and Hertz. We see that if the energy  $E$  of the electrons is less than the first excitation energy  $E_1 - E_0$ , the atoms remain in the ground state. If  $E$  becomes greater than  $E_1 - E_0$ , but remains less than  $E_2 - E_0$ , the atom can be brought by the collision into the first excited state, and consequently when it falls back into the ground state radiates only the line  $\nu_1 = \nu_{10}$ . If  $E + E_0$  lies between  $E_2$  and  $E_3$ , the atom which is struck can pass into either the first or the second excited state, and so can radiate the lines  $\nu_{20}$ ,  $\nu_{21}$ , and  $\nu_{10}$ ; and similarly in other cases (figs. 7a, 7b, Plate V). But we can also measure the energies of the electrons after the collision, by causing them to enter an opposing field of known potential difference, and observing the number of electrons which pass through it. In this way also the energy relation was found to be fulfilled exactly to this extent, that the energy loss of the electron due to the collision with an atom was just equal to an excitation energy  $E_n - E_0$  of the atom.



Ch. III, Fig. 9.—Transmutation of lithium on bombardment with protons. Pairs of  $\alpha$ -particles are shot out in opposite directions. (From *Proc. Roy. Soc. A*, Vol. 141.)



Ch. IV, Fig. 9.—Diffraction of electrons by thin silver foil. (After H. Mark and R. Wierl.) The velocity of the electrons (accelerating potential 36 kilovolts) corresponds to a de Broglie wave-length of  $0.0645 \text{ \AA}$ . (Exposure  $\frac{1}{10}$  sec.)



Ch. IV, Figs. 7a, 7b. Excitation of spectral lines by electronic collision (after Hertz). Only those lines appear in the spectrum whose excitation potential (the number in brackets) is smaller than the energy of the electrons (given under the spectra). Fig. 7a refers to mercury, 7b to helium. The wave-lengths are stated in Å.



#### 4. Compton Effect.

The phenomena described up to this point prove only that energy exchange between light and atoms, or between electrons and atoms, takes place by quanta. The *corpuscular nature* of light itself is proved in the most obvious way by the laws of *frequency change in the scattering of X-rays*. We have in an earlier chapter (Chap. III, p. 51) discussed the classical theory of the scattering of X-rays at comparatively weakly bound (nearly free) electrons, and reached the result that the scattered radiation has always the same frequency  $\nu$  as the primary radiation; for the electron vibrates in the same rhythm as the electric vector of the incident wave and, like every oscillating dipole, generates a secondary wave of equal frequency.

Compton (1922) investigated the scattering of X-rays by a block of paraffin, and found that the radiation scattered at an angle of less than  $90^\circ$  possesses a greater wave-length than the primary radiation, so that the  $\nu'$  of the scattered wave, contrary to the prediction of the classical theory, is smaller than the  $\nu$  of the

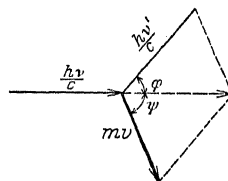


Fig. 8.—Compton Effect. A light quantum on colliding with an electron transfers part of its energy to the latter, and its wave-length becomes greater after the scattering.

incident radiation (see fig. 8). On the principles of the wave theory, this phenomenon is unintelligible.

The result, however, can be explained at once (Compton and Debye) if, taking the corpuscular point of view, we regard the process as one of elastic collision of two particles, the electron and the light quantum (fig. 8). For if the light quantum  $h\nu$  strikes an electron, it will communicate kinetic energy to the electron, and therefore will itself lose energy. The scattered light quantum will therefore have a smaller energy  $h\nu'$ . The exact calculation of the energy loss proceeds as in the case of the collision of two elastic spheres; the total momentum must be the same after the collision as before, likewise the total energy. The full calculation is given in Appendix X, p. 279; here we merely quote the result. The *Compton formula* for the change of wave-length of the light quantum due to the scattering process runs:

$$\Delta\lambda = 2\lambda_0 \sin^2 \frac{\phi}{2},$$

$$(\lambda_0 = \frac{h}{mc} = 0.0242 \text{ \AA.}, \text{ the Compton wave-length}).$$

The increase in the wave-length is accordingly independent of the wave-length itself, and depends only on  $\phi$ , the angle of scattering. The theory is susceptible of searching experimental tests, and is found to be thoroughly in accord with the facts. In the first place, Compton himself confirmed that the change of wave-length is correctly given by the Compton formula. The recoil electrons, which according to the theory are necessarily produced in the scattering process, and take over the energy loss  $h\nu - h\nu'$  of the light quanta in the form of kinetic energy, were successfully demonstrated by Compton; and that not only in scattering by solid bodies, but also in the Wilson chamber, where the tracks of the recoil electrons can be seen directly.

But, as has been shown by Compton and Simon, we can take a further step and test experimentally the relation between the angles of scattering  $\phi$  and  $\psi$  of the light quantum and the electron. Certainly a light quantum shows no track in the cloud chamber, but all the same we can determine the direction of the scattered quantum, provided it is scattered a second time and again liberates a recoil electron; the direction of the scattered quantum being found as that of the line joining the initial points of the tracks of the two recoil electrons. Although in the actual carrying out of the experiment there is a considerable amount of uncertainty, owing to the fact that several tracks may be present, and it is not always possible to determine a pair uniquely as corresponding to each other in the foregoing sense (i.e. produced by one and the same quantum), still Compton and Simon were able to establish agreement between theory and experiment with a fair amount of certainty.

Further confirmation of our ideas about the mechanism of the Compton effect was produced by Bothe and Geiger. They caused X-rays to be scattered in hydrogen, and with a Geiger counter recorded when recoil electrons made their appearance; by means of a second counter they determined the instants at which scattered light quanta appeared. (The counter, it is true, did not react directly to a light quantum striking it, but it did react indirectly to the electrons liberated by the light quantum in the chamber of the counter.) They succeeded in this way in establishing that the emission of the recoil electrons took place at the same moment as the scattering of the light quantum.

Recent investigations by Bothe, Jacobsen and others (1936) with powerful modern equipment (made, it is true, to disprove a negative result of Shankland), have definitely confirmed the simultaneous appearance of the recoiling light quantum and electron.



In the Compton scattering we have therefore a typical example of a process in which radiation behaves like a corpuscle of well-defined energy (and momentum); an explanation by the wave theory of the experimental results which we have described seems absolutely impossible. On the other hand, interference phenomena are quite irreconcilable with the corpuscular view of radiation. Until a few years ago, to explain this contradiction in the theory of light seemed to be beyond the bounds of possibility.

### 5. Wave Nature of Matter. De Broglie's Theory.

The dilemma was still further intensified, when in 1925 de Broglie propounded the hypothesis that the same dualism of wave and corpuscle as is present in light may also occur in matter. A material particle will have a matter wave corresponding to it, just as a light quantum has a light wave; and in fact the connexion between the two opposing aspects must again be given by the relation  $E = h\nu$ . Further, since from the standpoint of the theory of relativity energy and momentum are entities of the same kind (momentum is the spatial part of a relativistic four-vector, whose time component is energy), it is obviously suggested that for consistency we should write  $p = h\tau$ ; if  $\nu$  denotes the number of vibrations per unit time,  $\tau$  must signify the number of waves per unit length, and therefore be equal to the reciprocal of the wave-length  $\lambda$  of the wave motion; hence  $p = h/\lambda$ .

The extension of the wave idea from optics to mechanics can be carried through consistently. Before we go into this, however, we should like once more to point out the "irrationality" (as Bohr calls it) involved in thus connecting the corpuscular and wave conceptions:  $E$  and  $p$  refer to a point mass of vanishingly small dimensions;  $\nu$  and  $\tau$ , on the contrary, to a wave which is infinitely extended in space and time. The imagination can scarcely conceive two ideas which appear less capable of being united than these two, which the quantum theory proposes to bring into such close connexion. The solution of this paradox will occupy us later.

We shall first develop de Broglie's theory further from a formal point of view. A particle of momentum  $p$  in the direction of the  $x$ -axis and energy  $E$  is to be associated, then, with an infinitely extended wave of the form  $u(x, t) = Ae^{2\pi i(\nu t - \tau x)}$  by means of the two relations

$$E = h\nu, \quad p = h\tau.$$

This wave advances through space with a definite velocity, the *phase*

velocity  $u$ . The value of  $u$  we can find at once, by considering the (plane) surfaces of constant phase, viz.

$$\phi \equiv \nu t - \tau x = \text{const.}, \quad \text{or} \quad x = \frac{\nu}{\tau} t + \text{const.};$$

it follows that

$$u = \left( \frac{dx}{dt} \right)_{\phi=\text{const.}} = \frac{\nu}{\tau} = \nu \lambda.$$

Since  $\nu$  is in general a function of  $\lambda$ , and conversely, this equation embodies the law of *dispersion of the waves*.

But it must be remarked that the phase velocity is a purely artificial conception, inasmuch as it cannot be determined experimentally. In fact, to measure this velocity it would be necessary to affix a mark to the infinitely extended smooth wave, and to measure the velocity of the mark. But the only way in which we can make a mark on the wave train is by superimposing other wave trains upon it, which mutually reinforce each other at a definite place, and so create a hump in the smooth wave function. We have now to determine the velocity with which this hump moves; it is called the *group-velocity*.

A general method is given in Appendix XI (p. 281); here we confine ourselves to a simple special case, which gives the same result, and brings out the difference between phase-velocity and group-velocity with particular clearness. On the primary wave, which we suppose to have the form of  $u(x, t)$  above, we superimpose a wave with the same amplitude and a slightly different frequency  $\nu'$  and wave-length  $\lambda'$ . In this case, as we know, "beat" phenomena occur, and we make use of one of the beat maxima as the mark in our wave train. What we are interested in, then, is the velocity with which the beat maximum moves.

The superposition of the two wave trains gives us mathematically a vibration of the form

$$u(x, t) = e^{2\pi i(\nu t - \tau x)} + e^{2\pi i(\nu' t - \tau' x)}.$$

This expression can be written

$$\begin{aligned} u(x, t) &= e^{2\pi i\left(\frac{\nu+\nu'}{2}t - \frac{\tau+\tau'}{2}x\right)} \left\{ e^{2\pi i\left(\frac{\nu-\nu'}{2}t - \frac{\tau-\tau'}{2}x\right)} + e^{-2\pi i\left(\frac{\nu-\nu'}{2}t - \frac{\tau-\tau'}{2}x\right)} \right\} \\ &= 2 \cos 2\pi \left( \frac{\nu - \nu'}{2} t - \frac{\tau - \tau'}{2} x \right) e^{2\pi i\left(\frac{\nu+\nu'}{2}t - \frac{\tau+\tau'}{2}x\right)}. \end{aligned}$$

It therefore represents a vibration of frequency  $(\nu + \nu')/2$  and wave-

length  $2/(\tau + \tau')$ , the amplitude of which varies slowly (beats) relative to the vibration itself. The phase, as we deduce at once from the formula, moves with the velocity  $(\nu + \nu')/(\tau + \tau')$ . On the other hand, the maximum of the amplitude moves with the velocity  $(\nu - \nu')/(\tau - \tau')$ . In the limit when  $\nu'$  tends to  $\nu$ , and therefore  $\tau'$  to  $\tau$ , we find from this the value already found for the phase-velocity

$$u = \frac{\nu}{\tau} = \nu\lambda,$$

while the group-velocity is given by the limiting value

$$U = \lim_{\nu' \rightarrow \nu} \frac{\nu - \nu'}{\tau - \tau'}.$$

But this is simply, by definition, the derivative (differential coefficient) of the frequency  $\nu$  with reference to the wave-number  $\tau$ , if we regard  $\nu$  as a function of  $\tau$  (law of dispersion); hence we have

$$U = \frac{d\nu}{d\tau} = \frac{d\nu}{d(1/\lambda)}.$$

As is shown in Appendix XI (p. 281), this expression for the group-velocity holds perfectly generally.

We now apply these formulæ to the case of a *free particle* with velocity  $v$ . Writing  $\beta$  for  $v/c$ , and employing the relativistic formulæ for energy and momentum (see p. 271), we have here

$$\nu = \frac{E}{h} = \frac{m_0 c}{h} \frac{c}{\sqrt{1 - \beta^2}}, \quad \tau = \frac{p}{h} = \frac{m_0 c}{h} \frac{\beta}{\sqrt{1 - \beta^2}}.$$

The phase-velocity is given by

$$u = \frac{\nu}{\tau} = \frac{c}{\beta} = \frac{c^2}{v},$$

and is therefore greater than the velocity of light  $c$ , if the particle's velocity is less than  $c$ . The phases of the matter wave are therefore propagated with a velocity exceeding that of light—another indication that the phase-velocity has no physical significance. For the group-velocity we find

$$U = \frac{d\nu}{d\tau} = \frac{d\nu}{d\beta} \bigg/ \frac{d\tau}{d\beta},$$

it is exactly equal to the particle's velocity, for

$$\frac{dv}{d\beta} = \frac{m_0 c}{h} \frac{c\beta}{(1 - \beta^2)^{3/2}},$$

$$\frac{d\tau}{d\beta} = \frac{m_0 c}{h} \frac{1}{(1 - \beta^2)^{3/2}},$$

and therefore  $U = c\beta = v$ .

The relationship thus brought out is very attractive; in particular it tempts us to try to interpret a particle of matter as a wave packet due to the superposition of a number of wave trains. This tentative interpretation, however, comes up against insurmountable difficulties, since a wave packet of this kind is in general very soon dissipated. We need only consider the corresponding case in water waves. If we produce a wave crest at any point of an otherwise smooth surface, it is not long before it spreads out and disappears.

## 6. Experimental Demonstration of Matter Waves.

In view of the boldness of de Broglie's hypothesis, that matter is to be regarded as a wave process, the question of course at once suggested itself, whether and in what way the hypothesis could be put to the test of experiment. The first answer was given by Einstein (1925), who pointed out that the wave idea gives a simple explanation of the *degeneration of electrons in metals*, which expresses itself in the abnormal behaviour of metals in regard to their specific heat, and as an experimental fact was known to theoretical physicists before de Broglie. The subject will be discussed in detail in Chapter VII (§ 7, p. 217).

Further, it was known from the investigations of Davisson and Germer (1927) that in the reflection of beams of electrons by metals deviations occurred from the result to be expected on classical principles, more electrons being reflected in certain directions than in others, so that at certain angles a sort of selective reflection took place. The conjecture was first propounded by Elsasser (1925) that we have before us here a diffraction effect of electronic waves in the metallic lattice, similar to that which occurs in X-ray interference in crystals. The exact investigations which were then undertaken by Davisson and Germer actually gave interference phenomena in precisely the same form as the known Laue interference with X-rays.

Further experiments by G. P. Thomson, Rupp and others showed that when beams of electrons are made to pass through thin foils (metals, mica), diffraction phenomena are obtained, of the same kind

as the Debye-Scherrer rings in X-ray interferences (fig. 9, Plate V). Moreover, when the conditions of interference and the known lattice distances were used as data, it was found that de Broglie's relation between wave-length and the momentum of the electrons was completely confirmed.

The following rough calculation gives an indication of the kind of wave-lengths we have to deal with in beams of electrons. According to de Broglie, we have  $\lambda = h/p$  or, if we confine ourselves to electrons of not too high speed, so that we can leave relativistic corrections out of account,  $\lambda = h/mv$ . On the other hand, the velocity of the electrons is determined by the potential  $V$  applied to the cathode tube:  $\frac{1}{2}mv^2 = eV$ . Hence  $\lambda = h/\sqrt{(2meV)}$ , or, on inserting the numerical values ( $e = 4.77 \times 10^{-10}$  e.s.u.,  $m = 9.0 \times 10^{-28}$  gm.,  $h = 6.55 \times 10^{-27}$  erg sec.),

$$\lambda = \sqrt{\frac{150}{V}} \text{ \AA.},$$

when  $V$  is expressed in volts. Thus, to an accelerating potential of 10,000 volts there corresponds a wave-length of 0.122 \AA.; the wave-lengths of the electronic beams employed in practice therefore lie in approximately the same region as those of hard X-rays.

Although it is astonishing that the discovery of the diffraction of electrons was not made earlier, the fact must nevertheless be considered a piece of great good fortune for the development of atomic theory. What confusion there would have been if, soon after the discovery of cathode rays, experiments had been undertaken simultaneously on their charge and capability of deflection, and on their possibilities in regard to interference! Again, Bohr's theory of the atom, which later was to serve as the foundation of the expansion of atomic theory into wave mechanics, was essentially based on the assumption that the electron is an electrically charged corpuscle.

To-day the technique of electronic diffraction is so far advanced that it is already employed to some extent in industry instead of the earlier methods with X-rays for the purposes of research on materials. One advantage of using electrons is that decidedly higher intensities are available than there are with X-rays. Thus, for example, an interference photograph requiring an exposure of twenty-four hours with X-rays can be produced by means of electrons, with the same working data, in something like one second. Another advantage is that the wave-length of the beams of electrons can be varied at will by changing the tube potential; if the setting of the potentiometer is

changed, it can be seen at once on the screen how the whole diffraction image contracts or expands according as the wave-length is made shorter or longer.

Still more impressive than the establishment of the fact that electrons when they pass through a crystal lattice behave like waves, is the discovery of Stern and his collaborators (1932) that *molecular rays* (of  $H_2$  and He) also show diffraction phenomena when they are reflected at the surfaces of crystals. It was even possible to separate a beam of molecules of nearly uniform velocity with the help of a device similar to the arrangement for measuring the velocity of light: two toothed wheels rotating on the same axis. De Broglie's equation was confirmed for these particles with an accuracy of about 1 per cent. Here, surely, we are dealing with material particles, which must be regarded as the elementary constituents not only of gases (we recall our discussion of the subject of the kinetic theory of gases), but also of liquids and solids. If we intercept the molecular ray after its diffraction at the crystal lattice, and collect it in a receiving vessel, we find in the vessel a gas which has still the ordinary properties.

These diffraction experiments on whole atoms show that the wave structure is not a property peculiar to beams of electrons, but that there is a general principle in question; classical mechanics is replaced by a new *wave mechanics*. For, in the case of an atom, it is clearly the centroid of all its particles (nucleus and electrons), i.e. an abstract point, which satisfies the same wave laws as the individual free electron. Wave mechanics in its developed form does actually render an account of this.

## 7. The Contradiction between the Wave Theory and the Corpuscular Theory, and its Removal.

In the preceding sections we have had a series of facts brought before us which seem to indicate unequivocally that not only light, but also electrons and matter, behave in some cases like a wave process, in other cases like pure corpuscles. How are these contradictory aspects to be reconciled?

To begin with, Schrödinger attempted to interpret corpuscles, and particularly electrons, as *wave packets*. Although his formulæ are entirely correct, his interpretation cannot be maintained, since on the one hand, as we have already explained above, the wave packets must in course of time become dissipated, and on the other hand the description of the interaction of two electrons as a collision of two wave packets in ordinary three-dimensional space lands us in grave difficulties.

The interpretation generally accepted at present goes back to the present writer. According to this view, the whole course of events is determined by the laws of probability; to a state in space there corresponds a definite probability, which is given by the de Broglie wave associated with the state. A mechanical process is therefore accompanied by a wave process, the guiding wave, described by Schrödinger's equation, the significance of which is that it gives the probability of a definite course of the mechanical process. If, for example, the amplitude of the guiding wave is zero at a certain point in space, this means that the probability of finding the electron at this point is vanishingly small.

The physical justification for this hypothesis is derived from the consideration of scattering processes from the two points of view, the corpuscular and the undulatory. The problem of the scattering of light by small particles of dust or by molecules, from the standpoint of the classical wave theory, was worked out long ago. If the idea of light quanta is to be applied, we see at once that the number of incident light quanta must be put proportional to the intensity of the light at the place concerned, as calculated by the wave theory. This suggests that we should attempt (Born, 1926) to calculate the scattering of electrons by atoms, by means of wave mechanics. We think of an incident beam of electrons as having a de Broglie wave associated with it. When it passes over the atom this wave generates a secondary spherical wave; and analogy with optics suggests that a certain quadratic expression formed from the wave amplitude should be interpreted as the current strength, or as the number of scattered electrons. On carrying out the calculation (Wentzel, Gordon) it has been found that for scattering by a nucleus we get exactly Rutherford's formula (p. 53; Appendix IX, p. 276, and XX, p. 307). Many other scattering processes were afterwards subjected to calculation in this way, and the results found in good agreement with observation (Born, Bethe, Mott). These are the grounds for the conviction of the correctness of the principle of associating wave amplitude with number of particles (or probability).

In this picture the particles are regarded as independent of one another. If we take their mutual action into account, the pictorial view is to some extent lost again. We have then two possibilities. Either we use waves in spaces of more than three dimensions (with two interacting particles we would have  $2 \times 3 = 6$  co-ordinates), or we remain in three-dimensional space, but give up the simple picture of the wave amplitude as an ordinary physical magnitude, and replace

it by a purely abstract mathematical concept (the second quantisation of Dirac, Jordan) into which we cannot enter. Neither can we discuss the extensive formalism of the quantum theory which has arisen from this theory of scattering processes, and has been developed so far that every problem with physical meaning can in principle be solved by the theory (Appendix XXII, p. 312). What, then, is a problem with physical meaning? This is for us the really important question, for clearly enough the corpuscular and wave ideas cannot be fitted together in a homogeneous theoretical formalism, without giving up some fundamental principles of the classical theory. The unifying concept is that of probability; this is here much more closely interwoven with physical principles than in the older physics (e.g. the kinetic theory of gases, § 2, p 3; § 6, p. 9). The elucidation of these relationships we owe to Heisenberg and Bohr (1927). According to them we must ask ourselves what after all it means when we speak of the description of a process in terms of corpuscles or in terms of waves. Hitherto we have always spoken of waves and corpuscles as given facts, without giving any consideration at all to the question whether we are justified in assuming that such things actually exist. The position has some similarity to that which existed at the time the theory of relativity was brought forward. Before Einstein, no one ever hesitated to speak of the *simultaneous* occurrence of two events, or ever stopped to consider whether the assertion of the simultaneity of two events at different places can be established physically, or whether the concept of simultaneity has any meaning at all. In point of fact Einstein proved that this concept must be "relativized", since two events may be simultaneous in one system of reference, but take place at different times in another. In a similar way, according to Heisenberg, the concepts corpuscle and wave must also be subjected to close scrutiny. With the concept of corpuscle, the idea is necessarily bound up that the thing in question possesses a perfectly definite momentum, and that it is at a definite place at the time considered. But the question arises: can we actually determine exactly both the position and the velocity of the "particle" at a given moment? If we cannot do so—and as a matter of fact we cannot—i.e. if we can never actually determine more than one of the two properties (possession of a definite position and of a definite momentum), and if when one is determined we can make no assertion at all about the other property for the same moment, so far as our experiment goes, then we are not justified in concluding that the "thing" under examination can actually be described as a particle in the usual



sense of the term. We are equally unjustified in drawing this conclusion even if we can determine both properties simultaneously, if neither can then be determined exactly, that is to say, if from our experiment we can only infer that this "thing" is somewhere within a certain definite volume and is moving in some way with a velocity which lies within a certain definite interval. We shall show later by means of examples that the simultaneous determination of position and velocity is actually impossible, being inconsistent with quantum laws securely founded on experiment.

The ultimate origin of the difficulty lies in the fact (or philosophical principle) that we are compelled to use the words of common language when we wish to describe a phenomenon, not by logical or mathematical analysis, but by a picture appealing to the imagination. Common language has grown by everyday experience and can never surpass these limits. Classical physics has restricted itself to the use of concepts of this kind; by analysing visible motions it has developed two ways of representing them by elementary processes: moving particles and waves. There is no other way of giving a pictorial description of motions—we have to apply it even in the region of atomic processes, where classical physics breaks down.

Every process can be interpreted either in terms of corpuscles or in terms of waves, but on the other hand it is beyond our power to produce proof that it is actually corpuscles or waves with which we are dealing, for we cannot simultaneously determine all the other properties which are distinctive of a corpuscle or of a wave, as the case may be. We can therefore say that the wave and corpuscular descriptions are only to be regarded as complementary ways of viewing one and the same objective process, a process which only in definite limiting cases admits of complete pictorial interpretation. It is just the limited feasibility of measurements that defines the boundaries between our concepts of a particle and a wave. The corpuscular description means at bottom that we carry out the measurements with the object of getting exact information about momentum and energy relations (e.g. in the Compton effect), while experiments which amount to determinations of place and time we can always picture to ourselves in terms of the wave representation (e.g. passage of electrons through thin foils and observations of the deflected beam).

We shall now give the proof of the assertion that position and momentum (of an electron, for instance) cannot be exactly determined simultaneously. We illustrate this by the example of diffraction through a slit (fig. 10). If we propose to regard the passage of an electron

through a slit and the observation of the diffraction pattern as simultaneous measurement of position and momentum from the standpoint of the corpuscle concept, then the breadth of the slit gives the "uncertainty"  $\Delta x$ , in the specification of position perpendicular to the

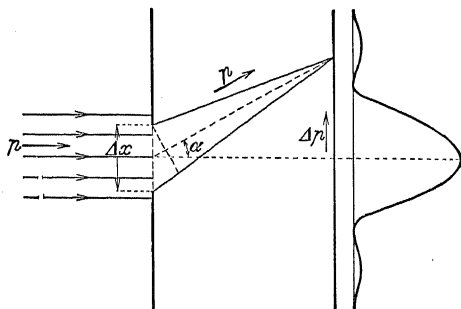


Fig. 10.—Diffraction of electrons at a slit

direction of flight. For the fact that a diffraction pattern appears merely allows us to assert that the electron has passed through the slit; at what place in the slit the passage took place remains quite indefinite. Again, from the standpoint of the corpuscular theory, the occurrence of the diffraction pattern on the screen must be understood in the sense that the individual electron suffers deflection at the slit, upwards or downwards. It acquires component momentum perpendicular to its original direction of flight, of amount  $\Delta p$  (the resultant momentum  $p$  remaining constant). The mean value of  $\Delta p$ , by fig. 11, is given by  $\Delta p \sim p \sin \alpha$ , if  $\alpha$  is the mean angle of deflection. We know that the experimental results can be explained satisfactorily on the basis of the wave representation, according to which  $\alpha$  is connected with the slit-width  $\Delta x$  and the wave-length  $\lambda = h/p$  by the equation  $\Delta x \sin \alpha \sim \lambda = h/p$ . Thus the mean added momentum in the direction parallel to the slit is given by the relation

$$\Delta p \sim p \lambda / \Delta x = h / \Delta x,$$

or

$$\Delta x \Delta p \sim h.$$

This relation is called *Heisenberg's uncertainty relation*. In our example therefore, it signifies that, as the result of the definition of the electron's position by means of the slit, which involves the uncertainty (or possible error)  $\Delta x$ , the particle acquires momentum parallel to the slit of the order of magnitude stated (i.e. with the indicated degree of uncertainty). Only subject to this uncertainty is its momentum known from the diffraction pattern. According to the uncertainty relation, therefore,  $h$  represents an absolute limit to the simultaneous measurement of co-ordinate and momentum, a limit which in the most favourable case we may get down to, but which we can never get beneath. In quantum mechanics, moreover, the uncertainty relation

holds generally for any arbitrary pair of "conjugated variables".

A second example of the uncertainty relation is the definition of position by a microscope (fig. 11). Here the order of ideas is as follows. If we wish to determine the position of an electron in the optical way by illuminating it and observing the scattered light, then it is clear, and known as a general rule in optics, that the wave-length of the light employed forms a lower limit to the resolution and accordingly to the exactness of the determination of position. If we wish to define the position as accurately as possible, we will employ light of the shortest possible wave-length ( $\gamma$ -rays). The employment of short-wave radiation implies, however, the occurrence of a Compton scattering process when the electron is irradiated, so that the electron experiences a recoil, which to a certain extent is indeterminate. We may investigate the circumstances mathematically. Let the electron under the microscope be irradiated in any direction with

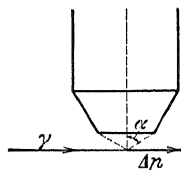


Fig. 11.—Determination of the position of an electron by means of the  $\gamma$ -ray microscope.

light of frequency  $\nu$ . Then by the rules of optics (resolving power of the microscope) its position can only be determined subject to the possible error

$$\Delta x \sim \frac{\lambda}{\sin \alpha},$$

where  $\alpha$  is the angular aperture. Now, according to the corpuscular view, the particle in the radiation process suffers a Compton recoil of the order of magnitude  $h\nu/c$ , the direction of which is undetermined to the same extent as is the direction in which the light quantum flies off after the process. Since the light quantum is actually observed in the microscope, this indeterminateness of direction is given by the angular aperture  $\alpha$ . The component momentum of the electron perpendicular to the axis of the microscope is therefore after the process undetermined to the extent  $\Delta p$ , where approximately

$$\Delta p \sim \frac{h\nu}{c} \sin \alpha.$$

Thus the order of magnitude relation

$$\Delta x \Delta p \sim h$$

holds good here also.

Just as every determination of position carries with it an uncertainty in the momentum, and every determination of time an

uncertainty in the energy (although we have not yet proved the latter statement), so the converse is also true. The more accurately we determine momentum and energy, the more latitude we introduce into the position of the particle and the time of an event. We give an example of this also, viz. the so-called *resonance fluorescence*. We have seen above (p. 73) that the atoms of a gas which is irradiated with light of frequency  $\nu_{10}$ , corresponding to the energy difference between the ground state and the first excited state, are raised to the latter state. They then fall back again to the lower state, at the same time emitting the frequency  $\nu_{10}$ ; and if the pressure of the gas is sufficiently low, so that the number of gas-kinetic collisions which occur while the atom remains in the excited state is negligible, then the whole energy which was absorbed will again be emitted. Thus the atom behaves like a classical resonator which is in resonance with the incident light wave, and we speak of resonance fluorescence.

But the energy of excitation of the atoms can also be utilized, not for re-emission of light, but for other actions, by introducing another gas as an indicator. If the latter consists, say, of not too rigidly bound diatomic molecules, the energy transferred in collisions with the excited atoms of the first gas can be utilized for dissociation (Franck, 1922). Again, if the added gas is monatomic, and has a lower excitation level than the first gas, it is itself caused to radiate by the collisions; this is called sensitized fluorescence (Franck). In any case we see that a fraction of the atoms of the first gas is certainly thrown into the excited state by the exciting light. We may take the following view of the matter. Excitation by monochromatic light means communication of exact quanta  $h\nu_{10}$  to the atom. We therefore know the energy of the excited atoms exactly. Consequently, by Heisenberg's relation  $\Delta E \Delta t \sim h$ , the time at which the absorption takes place must be absolutely indeterminate. We can satisfy ourselves that this is so, by considering that any experiment to determine the moment in question would necessarily require a mark in the original wave train—an interruption of the train, for example. But that means disturbing the monochromatic character of the light wave, and so contradicts the hypothesis. A rigorous discussion of the circumstances shows that, if the light is kept monochromatic, the moment at which the elementary act happens does actually altogether elude observation.

The uncertainty relation can also be deduced from the following *general idea*. If we propose to build up a wave packet, extending for a finite distance in the  $x$ -direction, from separate wave trains, we

need for the purpose a definite finite frequency-range in the monochromatic waves, i.e., since  $\lambda = h/p$ , a finite momentum-range in the particles. But it can be proved generally (Appendix XII) that the length of the wave packet is connected with the requisite range of momenta by the relation

$$\Delta p \Delta x \sim h.$$

The analogous relation

$$\Delta E \Delta t \sim h$$

can be derived in a similar way.

Bohr is in the habit of saying: the wave and corpuscular views are *complementary*. By this he means: if we prove the corpuscular character of an experiment, then it is impossible at the same time to prove its wave character, and conversely. Let us illustrate this further by an example.

Consider, say, *Young's interference experiment* with the two slits (p. 67); then we have on the screen a system of interference fringes. By replacing the screen by a photoelectric cell, we can demonstrate the corpuscular character of the light even in the fringes. It therefore appears as if we had here an experiment in which waves and particles are demonstrated simultaneously. Really, however, it is not so; for, to speak of a particle means nothing unless at least two points of its path can be specified experimentally; and similarly with a wave, unless at least two interference maxima are observed. If then we propose to carry out the "demonstration of a corpuscle", we must settle the question whether its path has gone through the upper or the lower of the two slits to the receiver. We therefore repeat the experiment, not only setting up a photoelectrically sensitive instrument as receiver, but also providing some contrivance which shows whether the light has passed through the upper slit (say a thin photographic film or the like). This contrivance in the slit, however, necessarily throws the light quantum out of its undisturbed path; the probability of getting it in the receiver (the screen) is therefore not the same as it was originally, i.e. the preliminary calculation by wave theory of the interference phenomenon is illusory. Thus, if pure interference is to be observed, we are necessarily precluded from making an observation of any point of the path of the light quantum before it strikes the screen.

We add in conclusion a few general remarks on the philosophical side of the question. In the first place it is clear that the dualism, wave-corpuscle, and the indeterminateness essentially involved therein,

compel us to abandon any attempt to set up a *deterministic theory*. The *law of causation*, according to which the course of events in an isolated system is completely determined by the state of the system at time  $t = 0$ , loses its validity, at any rate in the sense of classical physics. In reply to the question whether a law of causation still holds good in the new theory, two standpoints are possible. Either we may look upon processes from the pictorial side, holding fast to the wave-corpuscle picture—in this case the law of causation certainly ceases to hold; or, as is done in the further development of the theory, we describe the instantaneous state of the system by a (complex) quantity  $\psi$ , which satisfies a differential equation, and therefore changes with the time in a way which is completely determined by its form at time  $t = 0$ , so that its behaviour is rigorously causal. Since, however, physical significance is confined to the quantity  $|\psi|^2$  (the square of the amplitude), and to other similarly constructed quadratic expressions (matrix elements), which only partially define  $\psi$ , it follows that, even when the physically determinable quantities are completely known at time  $t = 0$ , the initial value of the  $\psi$ -function is necessarily not completely definable. This view of the matter is equivalent to the assertion that events happen indeed in a strictly causal way, but that we do not know the initial state exactly. In this sense the law of causation is therefore empty; physics is in the nature of the case indeterminate, and therefore the affair of statistics.

## CHAPTER V

# Atomic Structure and Spectral Lines

### 1. The Bohr Atom; Stationary Orbits for Simply Periodic Motions.

We have already had before us (Chap. IV, § 2, p. 69) a series of arguments to prove that the classical laws of motion cease to hold good in the interior of atoms. We recall in particular the existence of sharp spectral lines, and the great stability of atoms, phenomena which from the classical standpoint are perfectly unintelligible.

Bohr's explanation of spectra, which we expounded in § 3, p. 72, points out the road we must follow in setting up a new atomic mechanics. In fact, long even before the discovery of the wave nature of matter, an at least *provisional atomic mechanics* was successfully founded by Bohr and developed by himself and his collaborators, the most prominent of whom is Kramers.

The leading idea (Bohr's *correspondence principle*, 1923) may be stated broadly as follows. Judged by the test of experience, the laws of classical physics have brilliantly justified themselves in all processes of motion, macroscopic and microscopic, down to the motions of atoms as a whole (kinetic theory of matter). It must therefore be laid down, as an unconditionally necessary postulate, that the new mechanics, supposed still unknown, must in all these problems reach the same results as the classical mechanics. In other words, it must be demanded that, for the limiting case of large masses, or orbits of large dimensions, the new mechanics passes over into classical mechanics. We may obtain a concrete idea of the significance of the correspondence principle from the example of the hydrogen atom, which according to Rutherford consists of a massive nucleus, with an electron revolving round it. By the laws of classical mechanics (Kepler's first law) the orbit of the electron is an ellipse, or in special cases a circle; in the following discussion we shall confine ourselves to this special case. Let the radius of the circular orbit be  $a$ , and let it be described

with the angular velocity  $\omega$ . These two quantities are connected by the relation

$$a^3\omega^2 = \frac{Ze^2}{m},$$

which corresponds to Kepler's third law; it follows from the equality of the centrifugal force  $ma\omega^2$  and the Coulomb force of attraction  $Ze^2/a^2$ , where  $Z$  is the nuclear charge number (1 for H, 2 for  $\text{He}^+$ , 3 for  $\text{Li}^{++}$ , and so on).

What specially interests us here is the value of the energy of the revolving electron. By the principle of energy the sum of the kinetic and the potential energy is constant:

$$\frac{m}{2}a^2\omega^2 - \frac{Ze^2}{a} = E.$$

In this equation the energy, which is indeterminate to the extent of an additive constant, is so normalized that  $E$  denotes the work needed just to release the electron from its connexion to the atom, i.e. to bring it to rest ( $a\omega \rightarrow 0$ ) at an infinite distance ( $a \rightarrow \infty$ ) from the nucleus. By combining the two equations, we find

$$E = -\frac{m}{2}a^2\omega^2 = -\frac{Ze^2}{2a} = -\left(\frac{Z^2e^4m\omega^2}{8}\right)^{\frac{1}{3}}.$$

It follows from this that

$$\frac{|E|^3}{\omega^2} = \frac{Z^2e^4m}{8}$$

is constant. We may add that this equation also holds in the case when the orbit is an ellipse (with semi-major axis  $a$ ) if we understand that  $\omega$  means  $2\pi/T$ , where  $T$  is the period of a revolution (see also Appendix XIV, p. 286).

So much for the hydrogen atom, according to classical ideas and fundamental laws; to any orbital radius  $a$ , or to any angular velocity  $\omega$ , there corresponds a definite value  $E$  of the energy, while  $a$ , or  $\omega$ , can assume any value we choose.

In contrast with this we have the hypothesis of Bohr (§ 3, p. 72), according to which the atom can exist only in definite discrete states, and, at a transition from a state with the energy  $E_1$  to a state with the smaller energy  $E_2$ , emits the spectral line for which  $h\nu = E_1 - E_2$ . From the frequencies of the emission or absorption lines we can find the energies of the individual Bohr states.



Now the relations observed in the line spectrum of the hydrogen atom are known very exactly (fig. 1). It was Balmer (1885) who first showed that the lines situated in the visible region—all that were then known—can be represented by the formula

$$\nu = R \left( \frac{1}{2^2} - \frac{1}{m^2} \right),$$

with a spectral line corresponding to every integral value of  $m > 2$ ;  $R$  being a constant, the so-called Rydberg constant, the value of which was afterwards determined very exactly by Paschen (1916):

$$R = 109678 \text{ cm.}^{-1}.$$

Here  $\nu$ , in accordance with the usage of spectroscopists, stands for the *wave number*, i.e. the number of wave-lengths per cm.:  $\nu = 1/\lambda$ .

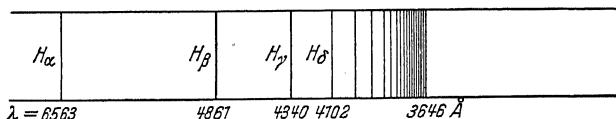


Fig. 1.—Diagram of the Balmer series, showing the notation used for the individual lines

The form of the above equation suggests the idea that the general law for the frequencies may be obtained by taking  $n^2$  instead of 4 ( $= 2^2$ ):

$$\nu = R \left( \frac{1}{n^2} - \frac{1}{m^2} \right).$$

In fact, on the basis of this formula the lines corresponding to the values  $n = 1, 3, 4, \dots$  could likewise be determined spectroscopically; they were not found in the earlier investigations, since they lie outside the visible spectral region. To-day the following series, called after their discoverers, are known:

- $n = 1$ : Lyman series (ultra-violet),
- $n = 2$ : Balmer series (visible),
- $n = 3$ : Paschen series (infra-red),
- $n = 4$ : Brackett series (infra-red),

. . . . .

It ought further to be particularly remarked that the agreement between the numbers given by the preceding formula and the values found spectroscopically is extraordinarily close. The accuracy of

spectroscopic work is high, and allows the position of the lines to be determined to 5 or 6 figures, yet the values calculated by the formula only differ from the observed values by at most a few units in the last place. The form of Balmer's equation—the difference of two expressions of the same type—suggests, in view of Bohr's theory, that these expressions should be brought into connexion with the energy levels (terms) of the Bohr atomic model:

$$E_n = -\frac{hR}{n^2} \text{ (Balmer term).}$$

The frequency emitted, when the atom makes a transition from a

state  $n$  to another state  $m$ , is then given, in agreement with the Balmer formula, by the relation

$$h\nu = E_n - E_m.$$

If we know the *term scheme* for an atom, we can at once read off from it the structure of the spectrum. For the hydrogen atom the term scheme has the form

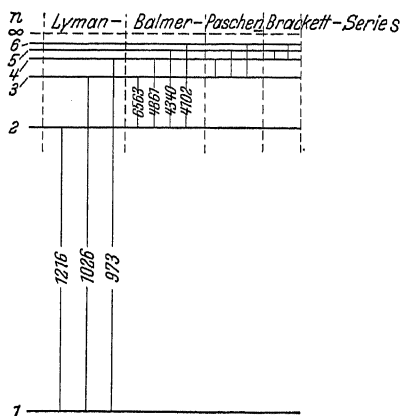


Fig. 2.—Term diagram for the hydrogen atom. The most important lines of the hydrogen spectrum (with some wave-lengths, in Å.) are shown as transitions between two terms.

shown above (fig. 2). The energy scale is normalized in such a way that its zero point corresponds to  $n \rightarrow \infty$ . Beginning at the lowest term  $E_1$ , the terms follow each other more and more closely, and asymptotically approach the limit  $E_\infty = 0$ .

In what precedes we have contrasted the statements of classical atomic mechanics with the results of experimental research. The latter, on Bohr's interpretation, give discrete energy levels with the values  $E_n = -hR/n^2$ , where  $R$  stands for a constant determined experimentally. It is of course the business of the new atomic mechanics to explain the value found for the Balmer term. The line of approach to the solution of the problem is indicated by the principle which we stated at the outset (p. 91)—the *correspondence principle*.

In the application to the hydrogen atom, this principle asserts that the higher the Bohr state of the atom is, the more closely the atom obeys the laws of classical mechanics; as  $n$  increases, the intervals

between the individual energy levels become smaller and smaller, the levels come closer and closer, and the atom approaches asymptotically the state of motion required by classical mechanics.

If then we calculate the emitted frequency by the Balmer formula, for the case when the initial and final states are highly excited states, we find, if  $m - n$  is small compared to  $m$  and  $n$ , the value

$$\nu \sim \frac{2R}{n^3} (m - n).$$

The lowest frequency emitted is got by taking  $m - n = 1$ , viz.

$$\nu_1 \sim \frac{2R}{n^3};$$

for  $m - n = 2$  we get a frequency twice as high, for  $m - n = 3$  one three times as high, and so on. The spectrum has therefore the same character as that of an electrically charged particle vibrating, on the classical theory, with the proper frequency  $\omega = 2\pi\nu_1$ , and the associated harmonic frequencies. If we now introduce the ground frequency  $\nu_1$  instead of the ordinal number  $n$  in the energy formula (Balmer term), we obtain (with  $\omega = 2\pi\nu_1$ ) the expression

$$E \sim -hR \left( \frac{\nu_1}{2R} \right)^{\frac{2}{3}} = - \left( \frac{Rh^3\nu_1^2}{4} \right)^{\frac{1}{3}} = - \left( \frac{Rh^3\omega^2}{16\pi^2} \right)^{\frac{1}{3}}.$$

This formula, in regard to the way in which it involves the frequency, has exactly the same structure as the formula which we obtained above (p. 92) for the energy of the classical revolving electron. As there, we have

$$\frac{|E|^3}{\omega^2} = \frac{Rh^3}{16\pi^2} = \text{constant}.$$

But, according to the correspondence principle, in the limiting case of very high values of  $n$ , i.e. of very low frequencies, the two formulæ must agree. It follows that

$$R = R_0 Z^2,$$

where

$$R_0 = \frac{2\pi^2 me^4}{h^3}.$$

Hence, on the basis of the correspondence principle, and under the assumption that the energy values are actually given by the Balmer terms, we have even at this early stage obtained an unambiguous

statement about the value of the Rydberg constant appearing in the Balmer term, and accordingly a possible means of testing the theory.

If we insert the values of  $e$ ,  $m$  and  $h$ , as determined by other methods, in the formula for  $R_0$ , we find  $R_0 \sim 3.290 \times 10^{15} \text{ sec.}^{-1}$ , or, when converted into wave numbers,  $R_0 \sim 109700 \text{ cm.}^{-1}$ ; a value which, within the limits of accuracy of our present knowledge of the values of  $e$ ,  $m$  and  $h$  (a few thousandth parts), agrees with the experimental value. Moreover, the formula depends in the right way on the ordinal number  $Z$ , as is shown by measurements on the ions  $\text{He}^+$ ,  $\text{Li}^{++}$ , . . . .

Additional evidence for the legitimacy of our argument is given by the fact that by a simple extension of the above formula the *motion of the nucleus* can also be taken into account, and that its effect, as thus found, upon the values of the terms has been completely verified by experiment. In the calculation given above, we proceeded as if the nucleus had infinite mass, and therefore assumed it to be at rest. In reality its mass, though certainly great compared to the mass  $m$  of an electron ( $\sim 1840 m$  for the hydrogen nucleus), is still finite. It follows that the nucleus must also move when the electron revolves, and that it describes a circular orbit round the common centre of mass with the same angular velocity  $\omega$  as the electron, since of course the centre of mass of the whole atom is to be at rest. If we denote the radii of the orbits of electron and nucleus by  $a_e$  and  $a_n$ , then by the definition of the centre of mass we have  $ma_e = Ma_n$ , where  $M$  is the mass of the nucleus. The equality of centrifugal force and Coulomb attraction gives then

$$ma_e\omega^2 = Ma_n\omega^2 = \frac{Ze^2}{(a_e + a_n)^2}.$$

In the formula for the energy the motion of the nucleus manifests itself, first, in the additional term  $\frac{1}{2}Ma_n^2\omega^2$  arising from the kinetic energy of the nucleus; and, secondly, in the fact that the potential energy contains, not as before the radius of the circular orbit of the electron, but the sum  $a_e + a_n$ . By the same process as before (p. 92), we obtain here the formula

$$E = - \left( \frac{Z^2 e^4 m M \omega^2}{8(m + M)} \right)^{\frac{1}{2}} = - \left( \frac{Z^2 e^4 m \omega^2}{8(1 + m/M)} \right)^{\frac{1}{2}}.$$

If now, using the correspondence idea, we equate this value of the energy to the value coming from the Balmer term for high

quantum numbers, we obtain for the Rydberg constant the expression:

$$R = \frac{R_0 Z^2}{1 + m/M}, \quad R_0 = \frac{2\pi^2 m e^4}{h^3}.$$

The motion of the nucleus manifests itself therefore in a correcting factor, which indeed amounts only to fractions of a thousandth part, but which the high accuracy of spectroscopic methods allows to be exactly tested.

Apart from the correcting factor, every energy term of even order of singly ionized helium coincides, according to the theory, with a hydrogen term, because  $Z = 2$  for  $\text{He}^+$ ; and therefore the second term of  $\text{He}^+$  is the same as the first of H, the fourth of  $\text{He}^+$  the same as the second of H, and so on. If, however, the motion of the nucleus is taken into account, these terms become separated, since the helium nucleus has four times the mass of the hydrogen nucleus, and therefore participates to a slighter degree in the motion of the electron than the latter. In fact, the formula for the energy terms of even order ( $2n$ ) of  $\text{He}^+$  gives the expressions

$$(E_{\text{He}^+})_{2n} = - \frac{4R_0 h}{1 + m/M_{\text{He}}} \frac{1}{4n^2} = - \frac{R_0 h}{(1 + m/M_{\text{He}})n^2},$$

while the hydrogen terms are given by

$$(E_{\text{H}})_n = - \frac{R_0 h}{(1 + m/M_{\text{H}})n^2}.$$

Every alternate helium line therefore almost, but not quite, coincides with a hydrogen line; the distance between them is easily calculated from the above formulæ; as to order of magnitude, it is approximately in the ratio  $m/M$  smaller than the wave-length of the lines themselves and therefore roughly of the order of magnitude of  $1 \text{ \AA}$ , so that it can easily be measured spectroscopically. Experimental tests of this result of the theory have verified it completely.

We may mention further that the value of the Rydberg constant, as measured by Paschen on the lines of hydrogen, which we have given above, is not equal to  $R_0$ , but equal to  $R_0/(1 + m/M_{\text{H}})$ . We thus find the value

$$R_0 = 109737 \text{ cm.}^{-1}.$$

It is at once obvious that the dependence of the frequency of the spectral lines on the mass of the nucleus will hold for other atoms also; it therefore becomes a possibility to discover isotopes spectro-

scopically. The most important case is that of the heavy hydrogen isotope D itself, for which our formula holds exactly, if we insert the nuclear mass  $M_D$ . In point of fact, as has already been mentioned (p. 60), the isotope D was first discovered by these spectroscopic methods.

The line of thought followed so far may be summarized as follows. Classical mechanics, on the basis of the picture of the electron revolving round the nucleus, certainly enables us to deduce formulæ for the connexion between orbital radius, frequency of revolution, and energy, but it is incapable of explaining the spectrum emitted by the atom. For the latter purpose we have, following Bohr, to introduce a new hypothesis, viz. that the atom only possesses certain definite energy levels  $E_n = -hR/n^2$ ; and it is the business of the new mechanics to explain these energy levels. From the correspondence principle, i.e. from the single requirement of asymptotic agreement of the new mechanics with the old for large quantum numbers, we have already been able to obtain definite information with regard to the connexion between the experimentally determined Rydberg constant and the atomic constants  $e$ ,  $m$  and  $h$ . But even so, there is still no explanation of why the Balmer terms occur at all; up to this point we only know that, if we make the special assumption that the formula  $E_n = -Rh/n^2$  is correct, then on account of the correspondence principle the factor of proportionality  $R$  must depend in a definite way on  $e$ ,  $m$  and  $h$ . It is for the new mechanics to explain these assumptions, or at least to make them intelligible. We may mention in anticipation that this was not accomplished until the introduction of wave mechanics. Nevertheless, we shall in the following pages deal briefly with the leading features and most important results of Bohr's quantum theory of the atom, since these are not only partly required as foundations for the structure of wave mechanics, but are also capable of explaining many experimental results.

In the first place, we collect here, in connexion with the preceding calculations, a few formulæ which are of great importance for the further development of Bohr's theory. If in the Bohr atom we think of the electrons as revolving in definite fixed orbits (ellipses, circles) round the nucleus, then with every energy level of the atom (in the special case of a circular orbit) we can associate a radius, an angular velocity, and so on, on the basis of the formulæ of classical mechanics. It then follows from the formulæ already given (p. 92) that (neglecting the motion of the nucleus) we have for the *radius*  $a$  of the circular orbit corresponding to the  $n$ th energy state

$$a = -\frac{Ze^2}{2E} = \frac{Ze^2n^2}{2hR_0Z^2} = a_0 \frac{n^2}{Z},$$

where

$$a_0 = \frac{h^2}{4\pi^2me^2} = 0.528 \text{ \AA}.$$

represents the radius of the first circular Bohr orbit in the hydrogen atom or, briefly, the "Bohr radius". Similarly, for the *angular velocity*, we find the expression

$$\omega = \frac{4\pi R_0 Z^2}{n^3} = \frac{\omega_0 Z^2}{n^3},$$

with

$$\omega_0 = 4\pi R_0 = 4.13 \times 10^{16} \text{ sec.}^{-1}.$$

Of special importance is the formula for the angular momentum (or, in other words, the moment of momentum) of the electron about the nucleus; from the two preceding formulæ it follows that

$$p = ma^2\omega = ma_0^2\omega_0n = n \frac{h}{2\pi}.$$

In the Bohr atom the angular momentum is equal to a multiple of  $h/2\pi$ . This fact is called the *quantum condition for angular momentum*.

Conversely, we may postulate this condition, and, by working backwards and using the formulæ of classical atomic mechanics, deduce the Balmer terms. Thus (p. 92), from

$$E = -\frac{Ze^2}{2a}, \quad \text{and} \quad a^3\omega^2 = \frac{Ze^2}{m},$$

it follows that

$$E = -\frac{Z^2e^4}{2ma^4\omega^2} = -\frac{Z^2e^4m}{2p^2},$$

where  $p = ma^2\omega$  denotes the angular momentum. If we substitute the value of  $p$  from the quantum condition, we find at once

$$E = -\frac{Z^2e^4m}{2} \frac{4\pi^2}{h^2n^2} = -\frac{hR_0Z^2}{n^2}.$$

It seems now a natural suggestion, that we should regard the quantisation condition for the angular momentum as an essential feature of the new mechanics. We therefore postulate that it is universally valid. At the same time, we must show by means of examples that the postulate leads to reasonable results. Although from the standpoint of

Bohr's theory the underlying reason for this quantisation rule remains entirely obscure, nevertheless in the further development of the theory it has justified itself by results.

As an example we take the case of the *rotating molecule*, which we treat as a rotator (i.e. as a rigid body which can rotate about a fixed axis). If  $A$  is its moment of inertia about this axis, then its kinetic energy when rotating with angular velocity  $\omega$  is, according to classical mechanics,

$$E = \frac{A}{2} \omega^2,$$

and there is no potential energy. The angular momentum about the axis is found from the energy by differentiating with respect to  $\omega$ :

$$p = A\omega.$$

We now apply the formula for the angular momentum, found for the case of the hydrogen atom:

$$p = n \frac{h}{2\pi},$$

where  $n$  denotes a whole number. Eliminating  $\omega$ , we obtain the energy formula:

$$E_n = \frac{p^2}{2A} = \frac{h^2}{8\pi^2 A} n^2.$$

This expression for the value of the energy is called the *Deslandres term*. To this term scheme there corresponds a definite emission spectrum. The transition from the  $n$ th state to the  $(n-1)$ th is bound up with the emission of a line of frequency

$$\nu = \frac{1}{h} (E_n - E_{n-1}) = \frac{h}{8\pi^2 A} (n^2 - (n-1)^2) = \frac{h}{4\pi^2 A} \left( n - \frac{1}{2} \right).$$

The spectrum consists of a sequence of equidistant lines, giving us the simplest type of *band spectrum*. Other transitions than those to the next lower (and next higher) state cannot occur, as follows from correspondence considerations; for, as we have already shown in the case of the hydrogen atom, and as we also see here at once, transition to the next quantum state but one, or to the next but two, for high quantum numbers involves the emission of double, or three times, the ground frequency, i.e. of harmonics. Now, classically, in the Fourier



expansion for circular motion (rotator), only the ground frequency occurs:

$$x = a \cos \omega t, \quad y = a \sin \omega t;$$

and the same of course holds good for the emitted radiation as calculated classically. Therefore, by the correspondence principle, in the quantum theory also no lines can be emitted which correspond to the classical harmonics in the Fourier resolution of the orbital motion; in other words, we have for the rotator the *selection rules*:  $n \rightarrow n - 1$  for emission transitions, and  $n \rightarrow n + 1$  for absorption transitions. We have only proved these rules, it is true, for high quantum numbers, since we have made use of the correspondence principle; but it is obviously suggested that we should postulate their validity for all quantum numbers.

## 2. Quantum Conditions for Simply and Multiply Periodic Motions.

The quantum condition for the angular momentum, which we have deduced from the Balmer term (which is given by experiment), and the application of which to the rotating molecule led us to the Deslandres term (which is in agreement with experiment), gives an indication of the way we have to follow in constructing the new mechanics. It comes to this clearly, that certain quantities can only take values which are whole numbers—they are called *quantisable quantities*—and the question arises of how in general these quantisable quantities are to be discovered. We shall try to make the matter clear with the aid of simple examples; and in Appendix XIII (p. 283) we give a general formulation of the problem on the basis of the Hamiltonian theory.

An idea of Ehrenfest's (1914) is helpful at this point. Consider a mechanical system, in which such quantisable variables occur, that is to say, quantities capable of integral values only. If we introduce a small disturbance, the new mechanics must hold good in the disturbed system just as in the undisturbed, so that the quantities in question must still have integral values. They must therefore at the beginning of the disturbance either have changed by a whole number at a jump, or have remained constant. If the actions affecting the system change slowly, the latter must be the case; when that is so, we say that these quantities are *adiabatically invariant*. As the converse of this argument, we may propound the theorem, that only adiabatically invariant quantities are quantisable. It is therefore necessary to discover such quantities, under the classical laws of motion.

We may show by a simple example how this can be done. Consider a *simple pendulum* (fig. 3) whose length can be altered, say by drawing the thread over a pulley. If we shorten the thread

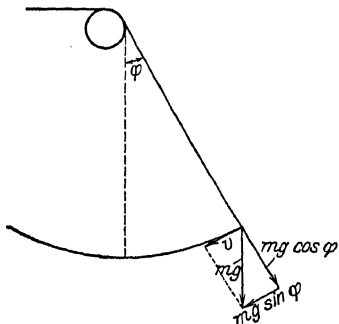


Fig. 3.—Simple pendulum of variable length. If the length is reduced slowly enough, the ratio of energy to frequency is constant.

slowly, we do work, first against gravity, secondly against the centrifugal force of the oscillating pendulum. Let the length of the thread be altered slowly from  $l$  to  $l + \Delta l$ ; to fix ideas we may suppose  $\Delta l$  to be a negative number, so that the pendulum is shortened. The component of the

weight, which stretches the thread, is  $mg \cos \phi$ , and the centrifugal force is  $ml\dot{\phi}^2$ , where  $\dot{\phi}$  is the angular velocity. The work done against gravity and centrifugal force is

$$A = \int_l^{l+\Delta l} (mg \cos \phi + ml\dot{\phi}^2) (-dl).$$

We now assume that the shortening of the thread takes place extremely slowly, so that during the process the pendulum oscillates to and fro a great many times. We can then disregard the variability of the amplitude of the oscillation with the length of the thread, and integrate over the motion on the supposition that the amplitude is constant. We thus obtain

$$A = -(mg \overline{\cos \phi} + m\overline{l\dot{\phi}^2})\Delta l,$$

where the bar indicates averaging over the undisturbed motion. If further we confine ourselves to small amplitudes, we can replace  $\cos \phi$  by  $1 - \phi^2/2$ . Thus

$$A = -mg\Delta l + (mg \frac{\overline{\phi^2}}{2} - m\overline{l\dot{\phi}^2})\Delta l = -mg\Delta l + \Delta W.$$

The first term corresponds to the elevation of the position of equilibrium, which does not interest us. The second part of the expression, i.e. the product of  $\Delta l$  by the expression in brackets, represents the increase  $\Delta W$  in the energy of the *pendulum motion*. Now the energy of the undisturbed pendulum motion is

$$W = \frac{m}{2} l^2 \dot{\phi}^2 + mgl(1 - \cos \phi),$$

where the first term represents the kinetic energy, and the second the potential energy relative to the position of rest. Replacing  $1 - \cos \phi$  by the approximation  $\phi^2/2$ , we get

$$W = \frac{m}{2} l^2 \dot{\phi}^2 + mgl \frac{\phi^2}{2}.$$

But this is just the energy function for a linear oscillator with the linear amplitude  $q = l\phi$ . The motion is therefore a simple harmonic vibration  $\phi = \phi_0 \cos \omega t$ , and we have therefore

$$\overline{\dot{\phi}^2} = \frac{\phi_0^2}{2}, \quad \overline{\phi^2} = \frac{\phi_0^2 \omega^2}{2},$$

from which, since  $\omega = 2\pi\nu = \sqrt{(g/l)}$ , it follows easily that

$$W = \frac{m l^2 \phi_0^2 \omega^2}{2} = mgl \frac{\phi_0^2}{2},$$

while, from the second term in the last expression for  $A$ ,

$$\Delta W = - \frac{m l \omega^2 \phi_0^2}{4} \Delta l = - \frac{W}{2l} \Delta l.$$

Hence we have

$$\frac{\Delta W}{W} = -\frac{1}{2} \frac{\Delta l}{l}.$$

But on the other hand, since  $\nu$  varies as  $l^{-\frac{1}{2}}$ , we have

$$\frac{\Delta \nu}{\nu} = -\frac{1}{2} \frac{\Delta l}{l},$$

so that

$$\frac{\Delta W}{W} = \frac{\Delta \nu}{\nu}.$$

This is a differential equation for  $W$  as a function of the frequency  $\nu$ , and its solution is

$$\frac{W}{\nu} = \text{const.} = J.$$

Thus, during the slow (adiabatic) shortening of the pendulum, this quantity  $J$  remains constant, so that without infringing Ehrenfest's principle mentioned above, we can put it equal to an integral multiple of  $h$ , i.e.

$$W = n h \nu.$$

We thus obtain the energy levels of the harmonic oscillator, in agreement with Planck's fundamental assumption.

The adiabatic invariants for other systems can be determined, theoretically, in a similar way. Still, this direct method is in general extremely troublesome, and one may well ask whether there is no simpler method of discovering invariants. We shall now show how it is possible to do so by means of a geometrical interpretation of the invariant quantity  $J = W/\nu$ , taking as a special case our previous example of the oscillator (simple pendulum with small oscillations).

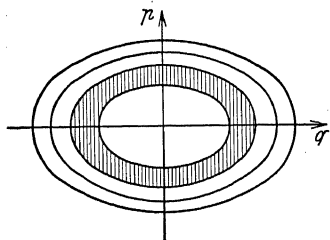


Fig. 4.—Phase paths for the linear oscillator; in the  $pq$ -plane the phase point describes an ellipse, whose area is an integral multiple of  $h$ .

We write down the energy once more, using a somewhat different notation ( $q = l\phi$ ,  $p = m\dot{q}$ ,  $f = mg/l$ ):

$$W = \frac{1}{2m} p^2 + \frac{f}{2} q^2.$$

In a  $pq$ -plane (fig. 4) this equation represents an ellipse with the semi-axes

$$a = \sqrt{2mW}, \quad b = \sqrt{2W/f},$$

as we see at once by writing the equation in the form

$$\frac{p^2}{2mW} + \frac{q^2}{2W/f} = 1.$$

The area of the ellipse is known to be

$$\oint p dq = ab\pi,$$

so that, in our case,

$$\oint p dq = 2\pi W \sqrt{(m/f)}.$$

(The symbol  $\oint$  means that we have to integrate over a whole period, i.e. here, over the whole circumference of the ellipse.) But we have  $2\pi\nu = \sqrt{(f/m)}$ , so that

$$\oint p dq = \frac{W}{\nu} = J.$$

The adiabatic invariant is therefore simply the area of the ellipse,

and the quantum postulate states that the area of the closed curve described in the  $pq$ -plane (phase plane), in one period of the motion, is an integral multiple of  $h$  (Debye, 1913).

The relation thus formulated is capable of immediate generalization. Consider in the first place, as an example with one degree of freedom, the case already treated above (p. 100), that of the *rotator*. Here the co-ordinate is the azimuth  $q = \phi$ , to which belongs, as canonically conjugated quantity, the angular momentum (or, in other words, the moment of momentum)  $p$ . In the free rotation  $p$  is constant, i.e. independent of the angle turned through. Thus

$$J = \oint p dq = p \oint dq.$$

If we represent the motion in the  $pq$ -plane, this integral is to be taken over the straight line  $p = \text{const.}$ , not over a closed curve. But we must observe that in this plane points with the same  $p$ , whose  $q$ -co-ordinates differ by  $2\pi$ , represent the same state of the rotator. Properly speaking, therefore, we should consider, not a  $pq$ -plane but a  $pq$ -cylinder (fig. 5) of circumference  $2\pi$ , so that the integral has now to be taken over the circumference of the cylinder, and has the

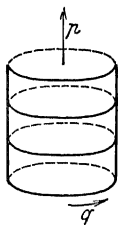


Fig. 5.—Representation of the phase path of a rotator on a cylindrical surface.

value  $2\pi$ . We therefore obtain  $J = 2\pi p$ . If we now assume the correctness of our quantisation rule

$$J = \oint p dq = nh,$$

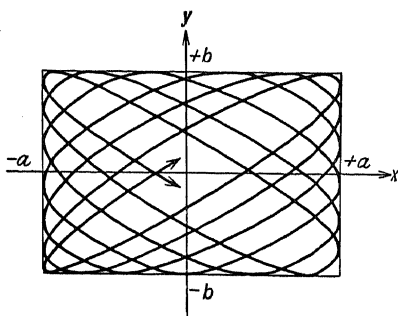
it follows that  $p = n(h/2\pi)$ , the formula which we found before in quite a different way for the hydrogen atom, and successfully applied to the rigid rotator for the interpretation of band spectra.

It is found that the quantisation rule

$$\oint p dq = nh$$

can be applied not only in systems with one degree of freedom, but also in complicated systems with many degrees of freedom, and always leads to results which agree with experience. The extension to more than one degree of freedom depends (Sommerfeld, Wilson, 1916) on the fact that in many cases co-ordinates  $q_1, q_2, \dots$  can be introduced such that the momenta associated with them have the property that  $p_1$  depends only on  $q_1$ ,  $p_2$  on  $q_2$ , and so on. Systems of this kind are

said to be *separable*. In general, they are multiply periodic. In that case the motion can be resolved into the superposition of simple periodic vibrations and their harmonics (so-called Lissajous' figures). As an example, consider the case of motion in a plane, where the rect-



angular co-ordinates  $x$  and  $y$  oscillate with two frequencies  $\nu_1$  and  $\nu_2$  (fig. 6). If  $\nu_1$  were equal to  $\nu_2$ , the path would be a circle, an ellipse or a straight line, according to the relation of the phases. If the ratio of  $\nu_1$  to  $\nu_2$

Fig. 6.—Curve representing the motion of a system with two degrees of freedom, in which the two frequencies  $\nu_1$  and  $\nu_2$  are incommensurable (Lissajous figure).

is a rational number, we again get closed orbits. If  $\nu_1$  and  $\nu_2$  are incommensurable, i.e. if their ratio is irrational, the curve does not close, but gradually fills the whole rectangle within which the variables range. For multiply periodic motions in general, the orbits are of this type. If, however, an orbital curve closes after a finite number of revolutions, there is in reality only one period, and there will be only one quantum condition of the type

$$\oint p dq = nh.$$

If the curve does not close, i.e. if two or more incommensurable periods are present, then there will be as many quantum conditions as there are periods:

$$\begin{aligned}\oint p_1 dq_1 &= n_1 h, \\ \oint p_2 dq_2 &= n_2 h, \\ &\dots\dots\dots\end{aligned}$$

This is the general case, also referred to as *non-degenerate*, whereas the case of coincident or commensurable periods is called a case of *degeneracy*. If  $u$  is the number of periods,  $v$  the number of degrees of freedom, then  $w = v - u$  is called the *degree of degeneracy*. In Appendix XIII (p. 283) we shall go into these relations a little more deeply.

As an example of the above rules we shall now give a discussion of the hydrogen atom, the complete quantisation of which was carried out by Sommerfeld. By Kepler's laws, the orbit of the electron round the nucleus is an ellipse; it is therefore simply periodic. Since the

electron has three degrees of freedom, this is a case of double degeneracy. In Appendix XIV (p. 286) we give the quantisation of the Kepler ellipse, which leads to the correct energy levels (Balmer terms).

The degeneracy is partly removed by taking into account the *relativistic variability of mass*, or dependence of the mass of the electron on its velocity. In this case the orbit, according to Sommerfeld, is given by a precessing ellipse (rosette); its major axis revolves in the plane of the ellipse round

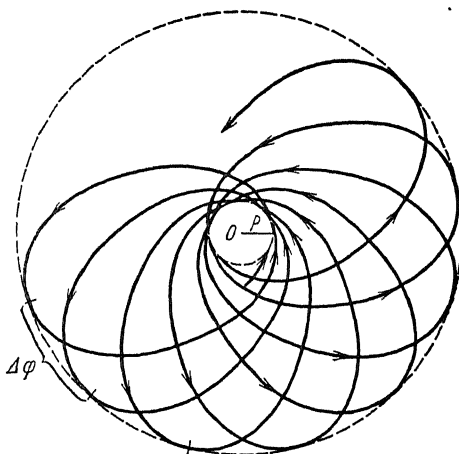


Fig. 7.—Orbit (rosette) of the electron about the nucleus, taking into account the relativistic variability of mass; the motion is doubly periodic, the perihelion being displaced by the angle  $\Delta\phi$  per revolution.

the nucleus with constant angular velocity (fig. 7). The orbit is now doubly periodic; besides the original period of revolution, which remains unchanged if the precession is slight, we have now the period of the precessional motion. In accordance with this, we have here *two* quantum conditions:

$$J_1 = nh, \quad J_2 = kh$$

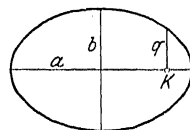


Fig. 8.—Elliptic orbit with the nucleus  $K$  as focus.

(compare Appendix XIV, p. 286);  $n$  determines the semi-major axis  $a$  of the approximate ellipse,  $k$  its semi-latus rectum  $q$ :  $a = na_0$ ,  $q = ka_0$  (fig. 8).

Further, the calculation shows that the total *angular momentum* is

$$p = k \frac{h}{2\pi},$$

and that the *energy* contains a term additional to the Balmer term:

$$E = -\frac{R_0 h Z^2}{n^2} + \epsilon(n, k).$$

From the first formula it follows that for  $k = 0$  the angular momentum vanishes; this gives the "pendulum orbits", in which the orbital ellipse degenerates into a straight line. For  $k = n$  we obtain the

greatest possible angular momentum for a fixed  $n$ ; this gives the circular orbits. For  $k < n$  we get the elliptic orbits. In the spectra, no sign has been found of the energy terms corresponding to the pendulum orbits; the terms not being realized, the theory must delete them from the term scheme. The reason given for their exclusion is that the pendulum orbits pass through the nucleus, so that the electron would collide with the nucleus, which of course is impossible.

The supplementary term  $\epsilon(n, k)$ , the value of which will be given later, expresses the minuter detail of the hydrogen lines; its effect is to split up every Balmer term into a number of energy terms, corresponding to the quantum number  $k$ . The spectral lines themselves are therefore made up of a system of finer lines, which are determined by the transitions between the energy levels of the upper state ( $n = n_1$ ,  $k = 1, 2, \dots, n_1$ ) and the lower state ( $n = n_2$ ,  $k = 1, 2, \dots, n_2$ ).

This is called the *fine structure* of the spectral lines. Its theory was given by Sommerfeld for the case of atoms of the hydrogen type ( $H$ ,  $He^+$ ,  $Li^{++}$ ), and was tested by Fowler and Paschen on the spectrum of singly ionized helium ( $He^+$ ), which was found in complete agreement with the theory. The test is easier with  $He^+$  than with  $H$  for this reason, that the energy terms of  $He^+$  are four times as far apart, on account of the nuclear charge number  $Z$  being doubled, whereas the

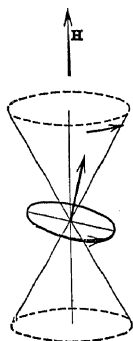


Fig. 9.—Top-like motion of the normals to the plane of an orbit, in an external field; the normal moves in a cone as shown.

corresponding factor for the fine structure splitting, as the theory shows, is 16; the fine structure of the lines of  $He^+$  is therefore more easily demonstrated and measured.

When the relativistic variability of mass is taken into consideration, the degeneracy of the hydrogen atom is certainly removed in part, but the motion is still simply degenerate. This degeneracy is connected with the fact that in the absence of an external field the orbital plane of the electron is fixed in space, and its orientation perfectly arbitrary. The degeneracy is only removed by the introduction of an external field. If the atom is brought into a homogeneous magnetic field  $H$ , a motion of precession of the orbital plane sets in, round the direction of the field (fig. 9). For the electron revolving in its orbital plane gives the atom angular momentum and also magnetic moment; considered as vectors, these are both at right angles to the orbital plane. The magnetic field takes hold of the magnetic moment and



seeks to turn it round into the field direction. The effort to do this is opposed by the rotational inertia due to the revolution of the electron, in a way known, sufficiently well for our purpose, from the theory of the top; the effect, as in the case of the mechanical top, is precessional motion.

This motion gives the third of the periodicities possible in the hydrogen atom. In accordance with our rules, it is to be quantised in the same way as the angular momentum of the orbital motion. This leads to the formula

$$p_{\phi} = m \frac{h}{2\pi} \quad (m = -k, -k + 1, \dots, +k).$$

Here  $p_{\phi}$  is the component of the angular momentum  $p$  in the direction of the field, so that this is quantised. Consequently the resultant angular momentum, the magnitude  $p$  of which must be an integral multiple of  $h/2\pi$ , can have only a finite number of possible inclinations to the direction of  $\mathbf{H}$ , in fact  $2k + 1$  inclinations (fig. 10), corresponding to the  $2k + 1$  possible values for its component  $p_{\phi}$

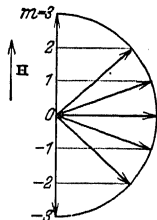


Fig. 10.—Example of quantisation of direction for the case  $k = 3$ ; the total angular momentum must be inclined at such an angle to the field direction that its projection  $m$  on that direction is a whole number (of units  $h/2\pi$ ).

parallel to  $\mathbf{H}$ . We speak of *quantising of direction*. On the experimental side, this quantising of direction can be demonstrated by the experiments of Stern and Gerlach, of which we shall give an account later (p. 167).

The precessional motion in the magnetic field contributes an additional quantity of energy  $m\nu_L h$ , where  $\nu_L = eH/(4\pi\mu c)$  is the so-called *Larmor frequency*,  $\mu$  being the mass of the electron. This value for the energy is obtained as follows. The potential energy of a magnet of moment  $M$  in a homogeneous magnetic field  $\mathbf{H}$ , with which it makes the angle  $\theta$ , is equal to  $(-)HM \cos \theta$ . In our case the angle  $\theta$  is defined by the quantising of direction:  $\cos \theta = p_{\phi}/p = mh/(2\pi p)$ . On the other hand, a revolving charge  $(-e)$ , which on account of its mass  $\mu$  gives rise to an orbital angular momentum  $p$ , produces by the rules of electron theory a magnetic moment of  $(-ep)/(2\mu c)$ . By combining these results, we find for the additional energy the value stated.

Every term is thus split up into  $2k + 1$  terms by the magnetic field. It is therefore to be expected that the spectral lines also possess the property of splitting in the magnetic field. That this is actually

so was found experimentally by Zeeman (1896), at a time when nothing was yet known of quantum theory, and attempted explanations of optical processes relied on classical mechanics. Zeeman's earliest observations showed that the lines in question, when observed transversally (magnetic field at right angles to the light path), are split up into three, the middle line being polarized parallel to the field, and the outside lines at right angles to the field; but that, when they are observed longitudinally (field parallel to light path), they split into two, which are circularly polarized in opposite directions. This case is now called the *normal Zeeman effect*. It was shown by Lorentz that a classical oscillator in a magnetic field may actually be expected to show these phenomena. Suppose the linear vibration of the oscillator resolved into two opposite circular motions at right angles to the field, along with a linear motion parallel to the field; then, by the action of the field, according to Lorentz's theory, the frequency of one of the circular motions is increased exactly by the Larmor frequency  $\nu_L$ , and that of the other circular motion is diminished by the same amount. This agrees with the splitting and polarization observed in one direction or another as described above.

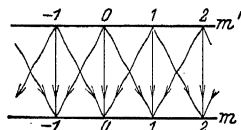


Fig. 11.—Transitions in a magnetic field; on account of the selection rule  $\Delta m = 0, \pm 1$ , only the transitions indicated by arrows occur; since the terms corresponding to the various values of  $m$  are equally far apart in a homogeneous magnetic field, each spectral line splits up in the magnetic field into three lines only (normal Zeeman effect).

It is now easy to see that our quantum theory leads to precisely the same result. In spite of the splitting into  $2k + 1$  terms, we obtain two, or three, component lines, according as we observe in the direction of the magnetic field, or perpendicular to it. In the light of the correspondence principle, this result can be understood at once. For the classical explanation of the Zeeman effect given by Lorentz implies that the atomic system, by the action of the magnetic field, is given an additional motion of rotation with the Larmor frequency  $\nu_L$ , that is to say, an additional pure circular motion without harmonics. By the correspondence principle this must also hold asymptotically for the treatment by quantum theory; hence, as in the rotator, we are led to the selection rules  $\Delta m = \pm 1$ . On the other hand, the component of the motion parallel to the field is not changed in this process; we therefore get the additional selection rule  $\Delta m = 0$ . For the possible transitions we therefore obtain the scheme of fig. 11. The splitting of the higher and lower terms is shown on horizontal lines, and the arrows correspond to the possible transitions. The height of

the energy jump is the same for all the arrows directed obliquely to the left, so that these yield one and the same radiated frequency; similarly with the arrows pointing vertically downwards, or towards the right. We therefore obtain a simple triplet (or doublet). Bohr's theory gives only the *normal Zeeman effect*. In reality (for not too strong fields), the splitting pattern consists in general, not of three lines only, but of a definitely greater number. This *anomalous Zeeman effect* cannot be understood without going beyond the concepts we have developed so far. We shall go into it more deeply in next chapter.

We pass now to the *spectra of the alkali atoms* (fig. 12, Plate VI). They are considered to arise in this way: an electron, the so-called radiating electron, moves in the field due to the nucleus and the rest of the electrons, and by itself causes the spectrum. This view is justified in the first place by the fact that in the alkalies, according to experiment, one electron is much more loosely bound than the rest, so that this electron is chiefly responsible for the chemical behaviour of the alkalies; on the other hand, we shall find, in the discussion of the periodic system in next chapter, that the remaining  $Z - 1$  electrons form so-called closed shells, round which the odd electron, i.e. the radiating electron, revolves. The field in which it moves is centrally symmetric, so that the potential depends only on the distance from the nucleus; the Coulomb field of the nucleus is "screened" by the remaining  $Z - 1$  electrons, and it is just these deviations from the Coulomb field which causes the differences between the *spectra of the alkalies* and that of hydrogen.

Here also the radiating electron moves in a precessing ellipse; the pure ellipse occurs as the form of the orbit in a pure Coulomb field only, any deviation from which, such as that determined by the variability of mass in the hydrogen atom, implies precession. We quantise this motion as before, and so obtain two quantum numbers  $n$  and  $k$ . By general agreement the terms are denoted by a number and a letter. The number indicates the principal quantum number  $n$ . For the specification of the azimuthal quantum number  $k$  the following notation has established itself:

$$\begin{array}{ccccccc} k = & 1 & 2 & 3 & 4 & \dots, \\ & s & p & d & f & \dots \text{ term.} \end{array}$$

Thus, e.g.,  $4d$  denotes the term with  $n = 4$  and  $k = 3$ . Since the precessional motion is purely periodic, the correspondence principle leads as before to the selection rule  $\Delta k = \pm 1$ . The only transitions which occur are therefore those from an  $s$ -term to a  $p$ -term, from a

$p$ -term to an  $s$  or  $d$ -term, and so on. For the sake of general distinction, the most important series of spectral lines have been given the following names:

$$\begin{aligned} np &\rightarrow 1s && \text{principal series,} \\ ns &\rightarrow 2p && \text{sharp series,} \\ nd &\rightarrow 2p && \text{diffuse series,} \\ nf &\rightarrow 3d && \text{fundamental series,} \\ &\dots && \dots \end{aligned}$$

(The sharp, diffuse and fundamental series are occasionally called respectively the second subordinate, first subordinate and Bergmann

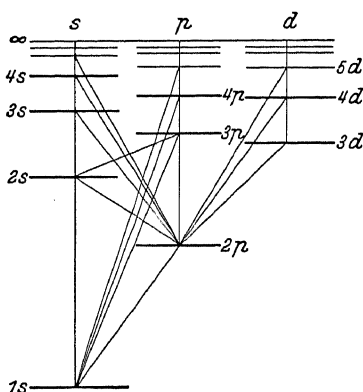
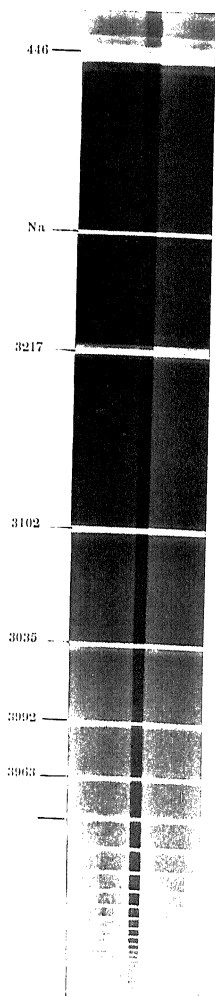


Fig. 13.—Term scheme of sodium (with not too high resolution). The transitions between the various levels give the emission lines.

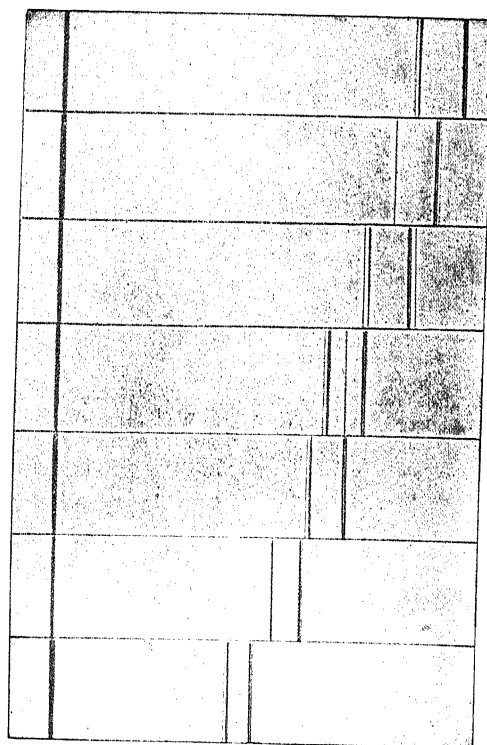
series.) In this list, we have on the left the terms between which the transition takes place, in the notation just explained; on the right, the names of the series obtained by giving different values to  $n$ . In the term scheme (fig. 13) it is to be understood that the energies increase upwards. For every azimuthal quantum number  $k$  there is a series of energy terms, which correspond to the various values of  $n$ , and which converge to a certain value (the *series limit*) when  $n \rightarrow \infty$ . The oblique lines represent the possible transitions; those to

the  $1s$ -term correspond to the spectral lines of the principal series, those from the  $s$ -series to the  $2p$ -term to the spectral lines of the sharp series, and so on. By the formation of such term schemes, founded on theoretical considerations, it has become possible to introduce order into the chaos of spectral lines as found experimentally.

As a last application of Bohr's theory we give a discussion of the origin of X-ray spectra. The essential feature distinguishing X-ray spectra from optical spectra is that the former are of the same type in all elements (fig. 14, Plate VI), whereas the latter, though of like structure for elements of like chemical character, are quite different for elements belonging to different columns of the periodic table. This means in the first place that X-ray spectra must have their source in the interior of the atom, whereas for optical spectra, just as for chemical behaviour, it is the outer region of atoms which matters.



Ch. V, Fig. 12. Emission spectrum of potassium vapour; a sodium line is shown between the two uppermost potassium lines.



Ch. V, Fig. 14. Rotating crystal photographs of the K-series for elements between arsenic ( $Z = 33$ ) and rhenium ( $Z = 47$ ). The figure shows the increase of hardness with increasing atomic number. (The marking on the extreme left is due to the undeflected primary beam.)



The lines found by experiment, which are denoted by  $K_\alpha, K_\beta, \dots, L_\alpha, L_\beta, \dots, M_\alpha, M_\beta, \dots$ , can be arranged in a term scheme, by taking advantage of the difference relations

$$K_\beta = K_\alpha + L_\alpha,$$

$$K_\gamma = K_\alpha + L_\beta = K_\beta + M_\alpha, \text{ \&c.},$$

which correspond to Ritz's combination principle; these have been tested experimentally with great exactness, and found to be fulfilled as stated.

The interpretation of X-ray spectra was given by Kossel (1917). In the atom the electrons are arranged in shells, there being a  $K$  shell, an  $L$  shell, &c. The electrons are most firmly bound in the  $K$  shell, less firmly in the  $L$  shell, still less firmly in the  $M$  shell, and so on. The energy levels indicated by horizontal lines in fig. 15 correspond to the electrons in the individual shells. The excitation of a  $K$ -line, according to Kossel, has to be pictured as follows. By some process (collision, absorption of light) an electron of the  $K$  shell is ejected from it. For this a definite minimum energy is requisite; all amounts of energy, which are greater than this least amount, can be absorbed. In absorption, we therefore find in the spectrum a sharp edge, the absorption edge; wave-lengths shorter than this are absorbed. Absorption lines, such as occur in the visible region, are not found in the X-ray region.

If an electron has been ejected in this way from the  $K$  shell, an electron of the  $L$  or  $M$  shell can fall down into the  $K$  shell, and a quantum of radiation is emitted; in this case the lines  $K_\alpha, K_\beta, \dots$  appear in the emission. A similar result follows in the case when the excitation causes an electron to be expelled from the  $L$  shell.

From his measurements Moseley was able to deduce the following law for the frequencies of the  $K_\alpha$ -lines of the various elements:

$$\nu = \frac{3}{4}R_0(Z - a)^2 = R_0(Z - a)^2 \left( \frac{1}{1^2} - \frac{1}{2^2} \right).$$

Here  $R_0$  is the Rydberg constant, which we have already had before us (p. 95);  $a$  is called the screening constant, and has approximately

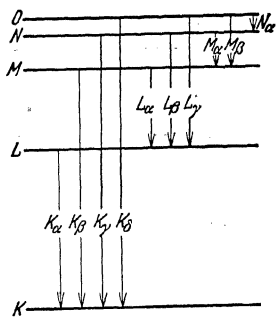


Fig. 15.—Term scheme of X-ray levels (after Kossel), with the transitions corresponding to the X-ray lines.

the same value 1.0 for all elements. The experimental facts, as stated in Moseley's law, imply therefore that in X-ray spectra we are concerned with terms of the hydrogen type ("hydrogen-like terms"), with the nuclear charge screened off. This result is to be understood in the sense that the innermost electrons, on account of the large nuclear charge, are bound almost solely to the nucleus, and therefore move in a similar way to the electron in the hydrogen atom, without being essentially disturbed in their orbits by the rest of the electrons

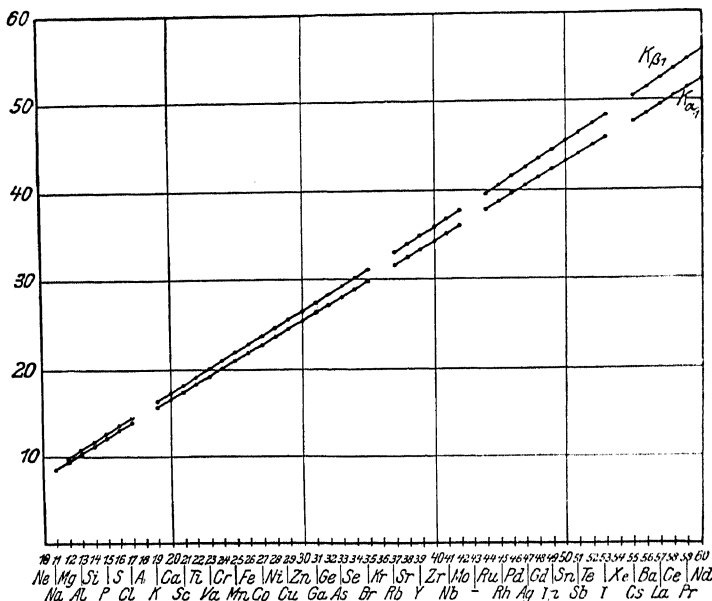


Fig. 16.—Diagram of points with abscissæ equal to the atomic number  $Z$ , and ordinates equal to the square roots of the frequencies of  $K$ -lines

in the atom. We therefore obtain a term scheme which corresponds to that of the alkali atoms. We must in fact associate with the electrons in the  $K$  shell the principal quantum number  $n = 1$ , with those of the  $L$  shell  $n = 2$ , and so on.

Further, an exact investigation shows that the  $K$ -level is single, the  $L$ -level triple, and the  $M$ -level quintuple, but the explanation of this is postponed to next chapter.

Moseley's law puts at our disposal a simple method of testing the order of succession in the periodic system. If we plot points in a diagram, with abscissa  $Z$ , and with ordinate the square root of the frequency of the  $K_{\alpha}$ -lines, the points of the diagram, as determined by



experiment, should lie on a straight line (fig. 16). If we gave an element a wrong place in the periodic table, i.e. associated the element with a wrong nuclear charge, then the corresponding point in the diagram would lie off the straight line. In this way it could be proved once again that, for instance, cobalt has a lower nuclear charge number (atomic number) than nickel, in spite of its greater mass; examples of this class are marked with a double arrow in the Table of p. 35. Another result of Moseley's work was the final determination of the gaps in the periodic table, i.e. of the elements which at that time were still unknown; in particular, the existence of hafnium ( $Z = 72$ ) could be predicted in this way before its real discovery (Coster and von Hevesy, 1923).

We may mention further that Bohr in his theory of the structure of the periodic system has substantially used only those simple properties which manifest themselves in the spectra of the alkalis and of X-rays.

Here we conclude our account of Bohr's theory. Although it has led to an enormous advance in our knowledge of the atom, and in particular of the laws of line spectra, it involves many difficulties of principle. At the very outset, the fundamental assumption of the validity of Bohr's frequency condition amounts to a direct and unexplained contradiction of the laws of the classical theory. Again, the purely formal quantisation rule, which stands at the head of the theory, is a foreign element which in the first instance is absolutely unintelligible from the physical point of view. We shall see later how both of these difficulties are removed in a perfectly natural way in wave mechanics.

Bohr's theory leaves some questions unexplained. Why must the pendulum orbits be excluded? The reason given, which points to the collision with the nucleus, is hardly cogent; moreover, it is outside the bounds of Bohr's theory. How does the anomalous Zeeman effect come about? The explanation of this, as we shall see (p. 140), requires us to use the fact that the electron possesses in itself mechanical angular momentum and magnetic moment. Finally, the calculation of the simplest problem of the type involving more than one body, viz. the helium problem, leads to difficulties, and to results contrary to experimental facts.

### 3. Matrix Mechanics.

The main reason for the break-down of Bohr's theory, according to Heisenberg (1925) is that it deals with quantities which entirely elude observation. Thus, the theory speaks of the orbit and the velocity

n electron round the nucleus, without regard to the consideration we cannot determine the position of the electron in the atom at without immediately breaking up the whole atom. In fact, in order to define its position with any exactness within the atom (whose diameter is of the order of magnitude of a few Ångström units), we must observe the atom with light of definitely smaller wave-length than this, i.e. we must irradiate it with extremely hard X-rays or with  $\gamma$ -rays; in that case, however, the Compton recoil of the electron is so great that its connexion with the atom is immediately severed, and the atom becomes ionized.

Thus, in Heisenberg's view, Bohr's theory fails because the fundamental ideas on which it is based (the orbit picture, the validity of the classical laws of motion, and so on) can never be put to the test. We move, therefore, in a region beyond experience, and ought not to be surprised if the theory, constructed as it is on a foundation of hypotheses which cannot be proved experimentally, partially fails in those deductions from it which can be subjected to the test of experiment.

If a logically consistent system of atomic mechanics is to be set up, no entities may be introduced into the theory except such as are physically observable—not, say, the orbit of an electron, but only the observable frequencies and intensities of the light emitted by the atom. Starting from this requirement, Heisenberg was able to lay down the leading principles of a theory, which was then developed by himself, Born and Jordan (1925), the so-called *matrix mechanics*, which is meant to replace the atomic mechanics of Bohr, and which in its applications has been brilliantly successful throughout. Although in form it is entirely different from the wave mechanics to be expounded later, yet in content, as Schrödinger has shown, the two are identical.

On this account, and especially in view of the complete novelty of the calculating technique, we do not consider matrix mechanics in any detail, but content ourselves with some brief indications. If we start from the frequencies

$$\nu_{nm} = \frac{E_n}{h} - \frac{E_m}{h},$$

as observable quantities, it is a natural suggestion that we arrange them in square array as follows:

$$\begin{pmatrix} \nu_{11} = 0 & \nu_{12} & \nu_{13} & \dots \\ \nu_{21} & \nu_{22} = 0 & \nu_{23} & \dots \\ \nu_{31} & \nu_{32} & \nu_{33} = 0 & \dots \\ \dots & \dots & \dots & \dots \end{pmatrix}.$$

If now we adhere to the arrangement in this scheme once for all, so that, say, the place in the fourth row and the second column is always associated with the transition from the fourth to the second quantum state, then we can also set out in similar square array the amplitudes  $\alpha_{nm}$  of the "virtual resonators" associated with the various frequencies emitted, where  $\alpha_{nm}^2$  then denotes the intensity of the frequency emitted:

$$\begin{pmatrix} \alpha_{11} & \alpha_{12} & \alpha_{13} & \dots \\ \alpha_{21} & \alpha_{22} & \alpha_{23} & \dots \\ \alpha_{31} & \alpha_{32} & \alpha_{33} & \dots \\ \dots & \dots & \dots & \dots \end{pmatrix}.$$

Similarly we can insert in an array of this kind other quantities connected with the transition  $n \rightarrow m$ .

The question now is: how do we calculate with these arrays? In the following remark of Heisenberg's is useful: if we multiply together two vibrational factors  $\alpha_{nk} = e^{2\pi i \nu_{nk} t}$  and  $\alpha_{km} = e^{2\pi i \nu_{km} t}$  by the  $\alpha_{nk} \alpha_{km}$ , then by Ritz's combination principle we get

$$e^{2\pi i \nu_{nk} t} e^{2\pi i \nu_{km} t} = e^{2\pi i (\nu_{nk} + \nu_{km}) t} = e^{2\pi i \nu_{nm} t},$$

which is a vibrational factor belonging to the same array, so that by the method specified for forming a product, we merely pass to another place in the square array, in conformity with the rule assigning places. We can proceed to define the product of two such arrays, in such a way that this product is again a square array of the same type. The multiplication rule, which Heisenberg deduced solely from experimental facts, runs:

$$(a_{nm})(b_{nm}) = \sum_k a_{nk} b_{km}.$$

It was remarked by Born and Jordan that this rule for multiplication is identical with one which has long been known in mathematics as the rule for forming the product of two "matrices", such as occur in the theory of linear transformations and the theory of determinants. We may therefore regard Heisenberg's square arrays as infinite matrices, and calculate with them by the known rules of the theory of matrices.

We now come to the central feature of matrix mechanics, which is this, that a representative matrix of the above type is associated with every physical magnitude. We can form a co-ordinate matrix, a momentum matrix, and so on, and then calculate with these matrices

in practically the very same way as we are accustomed to do with co-ordinates, momenta, and so on, in classical mechanics.

There is one essential distinction, however, between matrix and classical mechanics, viz. that when matrices are introduced as co-ordinates  $q_k$  and momenta  $p_k$ , their product is not commutative; it is no longer true, as it was in classical mechanics, that

$$p_k q_k - q_k p_k = 0.$$

The non-commutativity which presents itself here is not, however, of the most general kind, as the theory shows; for the left-hand expression, with a pair  $(p_k, q_k)$  of canonically conjugated variables, can take only the definite value

$$p_k q_k - q_k p_k = \frac{h}{2\pi i}.$$

These *commutation laws* (Born and Jordan, 1925) take here the place of the quantum conditions in Bohr's theory. The considerations by which their adoption is justified, as also the further development of matrix mechanics as a formal calculus, are for brevity omitted here. In the next section, however (§ 4, p. 121), it will be found that the analogous commutation laws in wave mechanics are mere matters of course. In Appendix XV (p. 291), taking the harmonic oscillator as an example, we show how and why they lead to the right result.

It may be mentioned in conclusion that the fundamental idea underlying Heisenberg's work has been worked out by Dirac in a very original way.

#### 4. Wave Mechanics.

Quite independently of the line of thought just explained, the problem of atomic structure has been attacked with the help of the ideas developed in the preceding chapter. According to the hypothesis of de Broglie (p. 79), to every corpuscle there corresponds a wave, the wave-length of which, in the case of rectilinear motion of the corpuscle, is connected with the momentum by the relation

$$\lambda = \frac{h}{p}.$$

It is only consistent to try to extend the theory by applying this wave idea to the atom, that is to say, to the electron revolving round the nucleus; in that case we have to picture the atom as a wave motion

round a particular point, the nucleus. What the theory has to do is to deduce the law of this motion.

As the first step, following de Broglie, we shall show that the quantum conditions of Bohr's theory can be interpreted at once on the basis of the wave picture. For this purpose we consider the simple case of a circular motion of the electron round a fixed point (fig. 17).

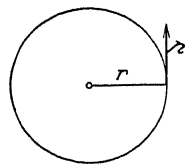


Fig. 17.—Motion in a circle from the corpuscular point of view; a particle moves in a definite orbit with the momentum  $p$ .

On Bohr's theory we have for this motion of revolution the quantum condition of angular momentum

$$p_\phi = rp = n \frac{h}{2\pi},$$

where  $p$  is the linear momentum  $mv$  of the electron. Now imagine a revolving wave instead of the revolving electron. If the radius of the circular orbit is very great, the same relation will hold for the revolving wave as for the plane wave, viz.

$$p = \frac{h}{\lambda}.$$

If we insert this value of the momentum in the preceding quantum condition of angular momentum, we obtain the equation

$$n\lambda = 2\pi r.$$

Here the term on the right represents the circumference of the circle: the formula states that this must be equal to a whole number of wave-lengths.

What then is the meaning of the quantum condition of angular momentum? If we try to construct

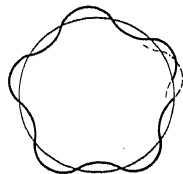


Fig. 18.—Motion of a wave in a circle; a single definite wave form is only possible when the circumference of the circle is a whole number of times the wave-length.

a wave motion along the circular orbit with an arbitrary wave-length  $\lambda$ , and for this purpose mark the wave train along the circumference (fig. 18), beginning at a point  $P$  with the "phase zero", then in general after a complete round we shall arrive at  $P$  with a phase differing from zero; after another round we again arrive with a new phase, and so on. We cannot *uniquely* associate a definite

phase with every point. It is otherwise in the case when we choose the wave-length so that

$$n\lambda = 2\pi r.$$

Here, in our construction, after a complete round we reach  $P$  again with the same phase as we started with. In this case the wave picture, or wave motion, is unique; a complete circuit changes nothing.

The quantum condition of angular momentum is therefore in this example identical with the requirement that the wave function of the corresponding vibratory process be one-valued. We therefore, as a general method, replace the hitherto unintelligible quantum condition of Bohr's theory by the obvious stipulation of *uniqueness (and finiteness) of the wave function* for the whole domain of the independent variables.

We pass now to the considerations which have led to the setting up of a differential equation—the *wave equation*—as the expression of the law of the wave motion in an atom. That the fundamental law takes the form of a differential equation is of course, by analogy with other vibratory processes, only to be expected. There is, naturally, no way of deducing the wave equation by strict logic; the formal steps which lead to it are merely matters of clever guessing.

We begin with the motion of a *free particle*; we describe the associated wave by a wave function

$$\psi = e^{2\pi i(\tau x - \nu t)} = e^{(2\pi i/h)(px - Et)}.$$

Here  $\nu$  and  $\tau$  are the frequency and wave number; according to de Broglie, they are connected with the energy and momentum by the equations

$$\nu = \frac{E}{h}, \quad \tau = \frac{1}{\lambda} = \frac{p}{h}.$$

By partial differentiation with respect to  $x$  and  $t$  we find

$$\frac{h}{2\pi i} \frac{\partial \psi}{\partial x} = p\psi, \quad -\frac{h}{2\pi i} \frac{\partial \psi}{\partial t} = E\psi.$$

We can now also read these equations in the converse way: the differential equations being given, it is required to find their solution. In the case when the particle moves in a *straight line*, and all values of  $x$  between  $-\infty$  and  $+\infty$  are permissible, the solution is represented by the function given above. If the particle moves along a circle of circumference  $l$ , and we denote by  $x$  the co-ordinate of a point, viz.

its distance from a fixed point on the circumference measured along the arc, then only values from 0 to  $l = 2\pi r$  are possible for  $x$ ; for increase of  $x$  by  $l$  brings us back to the same point. Since the value of  $\psi$  on the circle must be a single definite quantity, an increase of  $x$  by  $l = 2\pi r$  must make no change in the function. Now the general solution of the first of the two equations above is  $\psi = Ae^{(2\pi i/h)px}$ , which, when  $x$  is increased by  $l$ , is multiplied by  $e^{(2\pi i/h)pl}$ ; hence, if  $\psi$  is to be a "proper function", this factor must be equal to unity, i.e.

$$e^{(2\pi i/h)pl} = 1 = e^{2\pi in}$$

or

$$\frac{pl}{h} = n; \quad p = p_n = \frac{n\hbar}{l} = \frac{n\hbar}{2\pi r}.$$

This signifies that the preceding equation, in the case of circular motion, does not possess a permissible solution for all values of  $p$ , but only for the discrete "proper values"  $1\hbar/l, 2\hbar/l, 3\hbar/l, \dots$ .

The foregoing equations can also be interpreted as follows. When the wave function  $\psi$  is known, we obtain the corresponding momentum, or its  $x$ -component  $p_x$ , by differentiating the wave function partially as to  $x$ :  $(\hbar/2\pi i)\partial\psi/\partial x = p_x\psi$ . To the  $x$ -component of the momentum there belongs, as we say, the differential operator

$$p_x = \frac{\hbar}{2\pi i} \frac{\partial}{\partial x};$$

similarly for  $y$  and  $z$ . The operator belonging to the energy, on the other hand, is

$$E = -\frac{\hbar}{2\pi i} \frac{\partial}{\partial t}.$$

Operators, or entities which operate on any function, that is, which when applied to this function, generate another function, can be represented in the most diverse ways. Heisenberg's matrices are simply one definite kind of representation of such operators; another kind is the set of differential coefficients corresponding to the momentum components and the energy. In the latter kind of representation the Born-Jordan commutation laws admit of a simple interpretation; here, by what we have just seen,  $pq - qp$  simply means the application of the differential operator

$$\frac{\hbar}{2\pi i} \frac{\partial}{\partial q} q - q \frac{\hbar}{2\pi i} \frac{\partial}{\partial q}$$

to the wave function  $\psi$ . But

$$\frac{h}{2\pi i} \left( \frac{\partial q \psi}{\partial q} - q \frac{\partial \psi}{\partial q} \right) = \frac{h}{2\pi i} \psi.$$

Application of the operator  $pq - qp$  is therefore identical with multiplication of  $\psi$  by  $h/2\pi i$ , or, in symbolical form,  $pq - qp = h/2\pi i$ .

The formalism, which Schrödinger (1926) found suitable for the treatment of the wave theory of the atom, consists of the following rule. Write down the energy function  $H(p, q)$  of the Hamiltonian theory as an operator, by replacing  $p$  in it everywhere by  $(h/2\pi i)\partial/\partial q$ ; the operator corresponding to terms in  $p^2$  is obtained by repetition, viz.

$$p^2 = \frac{h}{2\pi i} \frac{\partial}{\partial q} \cdot \frac{h}{2\pi i} \frac{\partial}{\partial q} = -\frac{h^2}{4\pi^2} \frac{\partial^2}{\partial q^2}.$$

The energy operator  $H\left(\frac{h}{2\pi i} \frac{\partial}{\partial q}, q\right)$  is to be applied to a wave function  $\psi$ .

Instead of the energy equation  $H(p, q) - E = 0$ , we obtain the differential equation

$$\left\{ H\left(\frac{h}{2\pi i} \frac{\partial}{\partial q}, q\right) + \frac{h}{2\pi i} \frac{\partial}{\partial t} \right\} \psi = 0.$$

This is called *Schrödinger's equation* for the problem.

We have thus found the formalism, according to which any mechanical problem can be treated. What we have to do is to find the one-valued and finite solutions of the wave equation for the problem. If in particular we wish to find the *stationary solutions*, i.e. those in which the wave function consists of an amplitude function independent of the time and a factor periodic in the time (standing vibrations), we make the assumption that  $\psi$  involves the time only in the form of the factor  $e^{-(2\pi i/h)Et}$ . If we use this in Schrödinger's equation, we find an equation in which the time does not appear, viz.

$$\left\{ H\left(\frac{h}{2\pi i} \frac{\partial}{\partial q}, q\right) - E \right\} \psi = 0.$$

We have now before us a typical "*proper value*" problem: we have to find those values of the parameter  $E$  (the energy) for which this differential equation possesses a *solution which is one-valued and finite* in the whole domain of the variables (see Appendix XVIII, p. 300).

As an example of the method of forming the wave equation we take



first the *harmonic oscillator*. Its Hamiltonian function, already given above (p. 104), is

$$H = \frac{p^2}{2m} + \frac{f}{2} q^2.$$

From this, by the preceding rule, we obtain, as the Schrödinger equation

$$\left\{ -\frac{1}{2m} \frac{h^2}{4\pi^2} \frac{\partial^2}{\partial q^2} + \frac{f}{2} q^2 + \frac{h}{2\pi i} \frac{\partial}{\partial t} \right\} \psi = 0,$$

or, as the time-free equation of the stationary problem,

$$\left\{ \frac{h^2}{8\pi^2 m} \frac{\partial^2}{\partial q^2} + \left( E - \frac{f}{2} q^2 \right) \right\} \psi = 0.$$

More important for us is the case of the *hydrogen atom*. Here the Hamiltonian function is

$$H = \frac{1}{2m} (p_x^2 + p_y^2 + p_z^2) - \frac{e^2 Z}{r},$$

and the differential equation derived from it is

$$\left\{ \frac{1}{2m} \left( \frac{h}{2\pi i} \right)^2 \left( \frac{\partial^2}{\partial x^2} + \frac{\partial^2}{\partial y^2} + \frac{\partial^2}{\partial z^2} \right) - \frac{e^2 Z}{r} + \frac{h}{2\pi i} \frac{\partial}{\partial t} \right\} \psi = 0.$$

If we now introduce the usual differential symbol  $\Delta$  for Laplace's operator  $\partial^2/\partial x^2 + \partial^2/\partial y^2 + \partial^2/\partial z^2$ , and pass to the time-free equation by putting

$$\psi \sim e^{-(2\pi i/h)Et},$$

we obtain the wave equation

$$\left\{ \frac{h^2}{8\pi^2 m} \Delta + E + \frac{e^2 Z}{r} \right\} \psi = 0.$$

This is a wave equation in three-dimensional space, whose solutions we shall investigate later. It may make matters easier for the reader if we begin with corresponding problems in one and two dimensions; and for the sake of perspicuity we shall take our examples from classical mechanics (acoustics).

An example of this kind, in one dimension, is that of the *vibrating string*. Its differential equation follows from the theory of elasticity:

$$\frac{d^2 \psi}{dx^2} + \lambda \psi = 0.$$

Here  $\lambda$  depends on the mechanical conditions (tension, thickness of the string) and represents in effect the square of the frequency of vibration. (It may be remarked that in classical vibrational processes the proper value parameter always contains the square of the frequency of vibration, while in wave-mechanical problems the proper value parameter is given in general by the energy  $E = h\nu$ , and therefore contains the frequency in the first power.) The solutions of this differential equation are

$$\psi(x) = a \cdot \begin{cases} \cos\sqrt{\lambda}x, \\ \sin\sqrt{\lambda}x. \end{cases}$$

On account of the boundary conditions  $\psi(0) = 0$  and  $\psi(l) = 0$  (string of length  $l$  with ends fixed) the cosine vibration drops out at once as being inconsistent with the first condition. But even the sine vibration is not a solution of the boundary value problem, unless  $l\sqrt{\lambda}$  is an exact integral multiple of  $\pi$ , so that  $\psi$  vanishes when  $x = l$ . It is only for definite values of  $\lambda$  (the proper values) that we obtain possible forms of vibration; these are given in fact by

$$\psi(x) = \sin\left(\frac{n\pi x}{l}\right); \quad \lambda = \left(\frac{n\pi}{l}\right)^2.$$

The vibration with  $n = 1$  represents the fundamental, that with  $n = 2$  the first harmonic (octave), and so on. In the vibrational process, at definite points on the string there are *nodes*, that is to say, points which remain at rest throughout the vibration

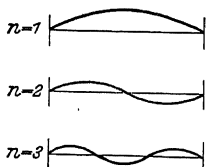


Fig. 19.—Vibrational forms of a string fixed at both ends. The fundamental ( $n = 1$ ), and the first two harmonics ( $n = 2$  and  $n = 3$ ).

(fig. 19). The number of nodes is determined by the parameter  $n$ , and in fact is clearly equal to  $n - 1$ .

As an actual example of such a system with one degree of freedom, in the quantum theory, we consider the *harmonic oscillator*, the wave equation for which we have already given above (p. 123). Its solution is dealt with in Appendix XVI (p. 295). Instead of Planck's energy levels  $E = nh\nu$ , wave mechanics, exactly like matrix mechanics (see Appendix XV, p. 291), gives the energy terms

$$E = (n + \frac{1}{2})h\nu.$$

The ground state (i.e. as regards energy, the lowest state,  $n = 0$ ) accordingly possesses a finite energy  $E = h\nu/2$  (zero-point energy). We shall

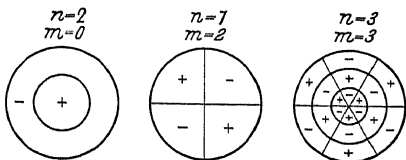
make use of this later (Chap. VIII, § 7, p. 255, App. XXXII, p. 343).

We next consider, as a two-dimensional example of a mechanical vibration, the vibrating *circular membrane*, the differential equation of which is

$$\left\{ \frac{\partial^2}{\partial x^2} + \frac{\partial^2}{\partial y^2} + \lambda \right\} \psi = 0.$$

Here also the proper value parameter  $\lambda$  depends on the nature of the plate and in effect represents the square of the frequency. The differential equation can be easily solved in polar co-ordinates (Appendix XVII, p. 297). In this case again we obtain possible forms of vibration

Fig. 20.—Some vibrational forms of a circular membrane fixed at the circumference; the number of radial nodal lines is here (in disagreement with the custom in wave mechanics) called  $n$ , that of the azimuthal nodal lines  $m$ ;  $n$  and  $m$  are the “quantum numbers” of the state of vibration.



for certain definite values of  $\lambda$  only. Instead of nodes, we have here nodal lines, of which indeed there are two kinds:

1. Lines for which  $r = \text{const.}$ ; these are defined by the radial ordinal number, or “quantum number”,  $n = 1, 2, \dots$
2. Lines for which  $\phi = \text{const.}$ , corresponding to the azimuthal ordinal number  $m = 0, 1, 2, \dots$

Fig. 20 shows several examples. The signs  $+$  and  $-$  in the different regions indicate that adjacent regions are always vibrating in opposite phases.

The solution can also be obtained in polar co-ordinates for the *hydrogen atom*, as a three-dimensional quantum problem; this is shown in Appendix XVIII, p. 298). In this problem, it should be added, we cannot speak of ordinary boundary conditions, since the domain over which the independent variables range is the whole of three-dimensional space. Instead of boundary conditions, we have now a rule with regard to the behaviour of the wave function at infinity. The natural condition to impose is that the wave function should vanish at infinity “more strongly” than  $1/r$ . This follows from the statistical interpretation of the square of the amplitude of the wave function, as the probability of the electron being found at a definite point of space. The condition is equivalent to this, that the electron must always be at a finite distance.

Taking this “boundary condition” into account, we obtain solutions of the wave equation which correspond to a bound electron

(elliptic orbit in the Bohr atom), only for definite discrete values of  $E$ . In this way we find as proper values exactly the Balmer terms with the correct Rydberg constant,  $E = -R\hbar/n^2$ . Here  $n$  is the principal quantum number. Besides this, the azimuthal quantum number  $l$  and the magnetic quantum number  $m$  also make their appearance. The number of nodal surfaces  $r = \text{const.}$  is  $n - l - 1$ ; for a given  $n$ ,  $l$  can therefore be any integer from 0 to  $n - 1$ ; as for  $m$ , it can take all values from  $-l$  to  $+l$ .

When the relativistic correction for the mass is taken into account, the energy depends on  $l$  also. Moreover, as follows from the form of this dependence,  $l + 1$  corresponds to the Bohr quantum number  $k$ , so that our nomenclature for the terms (see § 2, p. 111) has now to be understood as follows:

$$\begin{array}{ccccccc} l = & 0 & 1 & 2 & \dots & & \\ & s & p & d & \dots & \text{term.} & \end{array}$$

In a *magnetic field*,  $E$  depends on  $m$  also, indeed the extra term  $mv_z\hbar$  occurs as an addition to the energy, exactly as in Bohr's theory. Wave mechanics, so far as we have developed it up to the present, yields only the normal Zeeman effect (as above, § 2, p. 110). To the directional quantisation of Bohr's theory there corresponds here the finite number of values of  $m$ , i.e. of energy levels in the magnetic field; there are in fact  $2l + 1$  of these, in place of each term which occurs when there is no magnetic field. The splitting of the terms in an *electric field* (Stark effect) is correctly reproduced by wave mechanics, qualitatively and quantitatively.

The states of the hydrogen atom, as considered so far, correspond clearly to the elliptic orbits of the old Bohr theory; in both cases the electron remains at a finite distance. But in Bohr's theory there are also hyperbolic orbits; what corresponds to these in quantum mechanics? Clearly, solutions of the wave equation which do not disappear at infinity. In order to obtain them we must give up the boundary condition—vanishing of  $\psi$  at infinity—and look for solutions which at a great distance from the nucleus behave approximately like plane waves. In point of fact there are such solutions, and that for all positive values of the energy. Physically interpreted, they describe what happens when an electron coming from infinity passes near the nucleus and is deflected by it. It can actually be proved that Rutherford's scattering formula is strictly valid in wave mechanics also; we shall return to this again (§ 6, p. 130; Appendix XX, p. 304). Further, corresponding to the transitions between such states of posi-

tive energy and the quantum states of negative energy, we have now the processes of emission and absorption of light, which form a continuous spectrum, and are particularly important for the purpose of understanding the distribution of intensity in X-rays.

We add a few remarks on the wave mechanics of *many-body problems*. Here of course we are concerned with the solution of a wave equation in many-dimensional space; thus the calculation of the helium spectrum needs as many as six co-ordinates, and that of the lithium spectrum nine. It is clear that in these cases an exact solution is not to be looked for, so that we must be content with an approximate solution of the problem. The methods of a highly developed *perturbation theory* enable us to push this approximation as far as we please; the labour involved, however, increases without limit with the order of the approximation. The lowest terms of He,  $\text{Li}^+$  and Li have already been successfully calculated by this method, with results in good agreement with experiment (Hylleraas, 1930).

Deductions which are quite exact can be made from any *properties of symmetry* which the wave function must possess in virtue of the symmetrical character of the problem. The most important of these properties of symmetry is the one which is involved in the complete equivalence of the electrons, and their consequent interchangeability; the wave function must of course be the same, whether, say, the first electron is situated in the *K* shell and the second in the *L* shell, or the second in the *K* shell and the first in the *L* shell. This leads to general rules for the tabulation of the terms in atoms with several radiating electrons. Still, the results thus obtained are not immediately comparable with experiment, since in wave mechanics, so far as developed above, an essential principle is lacking, which was discovered by Pauli, and which will come before our notice in next chapter.

## 5. Angular Momentum (Moment of Momentum) in Wave Mechanics.

In Bohr's theory, angular momentum was of special importance for the classification of the spectral lines and the systematic arrangement of the terms; it was found that it corresponded to a quantum number  $k$ . This again led to the idea of the quantisation of direction, the experimental confirmation of which by Stern and Gerlach's experiment is perhaps the most impressive evidence we have of the fundamental difference between classical and quantum mechanics. The question is thus raised: how does wave mechanics deal with these things? Is it capable of taking account of quantisation of direction in a natural way?

In wave mechanics, angular momentum, just like linear momentum, has a differential operator corresponding to it, the components of the operator being

$$M_x = yp_z - zp_y = \frac{h}{2\pi i} \left( y \frac{\partial}{\partial z} - z \frac{\partial}{\partial y} \right),$$

$$M_y = zp_x - xp_z = \frac{h}{2\pi i} \left( z \frac{\partial}{\partial x} - x \frac{\partial}{\partial z} \right),$$

$$M_z = xp_y - yp_x = \frac{h}{2\pi i} \left( x \frac{\partial}{\partial y} - y \frac{\partial}{\partial x} \right).$$

What is the significance of these operators? The electronic state is given by the wave function  $\psi$ . In order to decide the question whether definite values of the components of angular momentum round the three co-ordinate axes belong to this state, we have to "apply" the above operators, according to the rules of wave mechanics, to the wave function  $\psi$ ; i.e. we perform the differentiations involved in the operators. There are now two possibilities: either this operation reproduces the wave function except for a constant factor, or it does not. In the first case, the wave function is also at the same time a proper function of the equation of angular momentum

$$M_x \psi = M_x \psi,$$

and the state represented by the proper function therefore possesses a definite angular momentum round the  $x$ -axis, whose value is given by the "proper value"  $M_x$ ; here  $M_x$  is an ordinary number (not, like  $M_x$ , an operator). If, however, when we apply the operator  $M_x$  to the wave function we obtain another function, which does not agree with the wave function except for a constant factor, that is to say, if  $\psi$  is not a proper function of the equation of angular momentum, this means that the electronic state in question is not associated with a fixed value of the angular momentum about this axis. In the case of the proper functions given in Appendix XVIII (p. 298), by our choice of the polar axis ( $z$ -axis) we have specially distinguished this axis from the beginning. According to Bohr's theory, the component angular momentum round this axis must be quantised. Now wave mechanics does in fact show that the proper values of the  $z$ -component of the angular momentum are integral multiples of  $h/2\pi$ . Thus, on introducing spherical polar co-ordinates

$$M_z = \frac{h}{2\pi i} \left( x \frac{\partial}{\partial y} - y \frac{\partial}{\partial x} \right) = \frac{h}{2\pi i} \frac{\partial}{\partial \phi},$$

and on applying this operator to the proper function of the state characterized by the quantum numbers  $n, l, m$ , we obtain at once on account of the way in which  $\phi$  is involved ( $e^{im\phi}$ )

$$M_z \psi_{nlm} = \frac{\hbar}{2\pi i} \frac{\partial}{\partial \phi} \psi_{nlm} = m \frac{\hbar}{2\pi} \psi_{nlm},$$

so that the proper values of the component angular momentum  $M_z$  are actually  $m\hbar/2\pi$ . In the other two co-ordinate directions, in the case before us, we may easily satisfy ourselves that we do not get proper values, and therefore do not get definite values of  $M_x$  and  $M_y$ . On the other hand, the value of the resultant angular momentum for an atom which can rotate freely (in the absence of an external field) is in all circumstances quantised. For brevity we refer to Appendix XIX (p. 302) for the requisite calculations, and give only the result here: the square of the value of the resultant

$$M^2 = M_x^2 + M_y^2 + M_z^2$$

is, according to wave mechanics,  $l(l+1)(\hbar/2\pi)^2$ , not, as in Bohr's theory,  $l^2(\hbar/2\pi)^2$ . This special feature is characteristic of the whole wave mechanics of the atom, and in next chapter we shall often meet it. Moreover, as we shall also see there, the splitting patterns of the terms in the anomalous Zeeman effect give direct confirmation of the fact that the square of the angular momentum is actually equal to  $l(l+1)$ , and not to  $l^2$ . If, however, we disregard this difference, we can apply in wave mechanics the representation of the angular momentum by a vector diagram, known from Bohr's theory; we therefore here also represent the total angular momentum by a vector  $l$ , with regard to which we must note once for all that it has the absolute value  $\sqrt{l(l+1)}$ . In the case when the  $z$ -direction is specially marked out, say by the fact that an (extremely) weak magnetic field  $H$  acts along it, then the component in this direction of this vector angular momentum is capable of integral values  $m$  only (in units  $\hbar/2\pi$ ), and this state of matters continues to hold even in the limiting case  $H \rightarrow 0$ . Here, and especially in next chapter, for the sake of greater perspicuity we take over the vector representation of the angular momentum precessing round the specially distinguished axis, and having a component in the direction of the axis which can take only integral values.

We expressly emphasize, however, that this idea cannot in this way without more ado be brought into harmony with the conceptual structure of wave mechanics. For the latter purpose, and particularly

for the proof that two angular momenta may be compounded vectorially in wave mechanics in the same way as in Bohr's semi-classical theory, higher mathematical methods are required, especially so-called group theory. For this reason we cannot go further into these questions at this point. It may be mentioned, however, that further development has led to the electron being regarded, not as a particle, defined by three space co-ordinates, but as a top-like structure, possessing an angular momentum of its own. This "spin" theory of the electron will be dealt with later (Chap. VI, p. 136).

In conclusion, we have still to consider the *meaning of the wave function* itself; so far, we have obtained it as a mere by-product, so to speak, in the search for proper values. But in a vibrational process knowledge of the amplitude is at least as important as knowledge of the proper frequency; similarly, it is to be expected that in wave mechanics great physical significance attaches to the wave function  $\psi$ , or rather, to the square of its modulus, since of course the instantaneous value of the oscillating function itself cannot play any part, on account of the high frequency. The reason for taking the square of the modulus is that the wave function itself (because of the imaginary coefficient of the time derivative in the differential equation) is a complex quantity, while quantities susceptible of physical interpretation must of course be purely real.

## 6. The Statistical Interpretation of Wave Mechanics.

We have already mentioned the interpretation of the wave function given by the author (p. 83). Let the proper function corresponding to any state be  $\psi_x$ ; then  $|\psi_x|^2 dv$  is the probability that the electron (regarded as a corpuscle) is in the volume element  $dv$ .

This interpretation is almost self-evident, if we consider, not the quantum states proper (with discrete, negative energy-values), but the states of positive energy, which correspond to the hyperbolic orbits of Bohr's theory. We have then to solve a wave equation

$$\left\{ \frac{\hbar^2}{8\pi^2m} \Delta + E - V(r) \right\} \psi = 0,$$

where, instead of the Coulomb potential  $-e^2Z/r$ ,  $V(r)$  is written somewhat more generally, in order to take into account possible screening of the action of the nucleus by firmly bound electrons. For particles entering the atom with very high speeds, and therefore very large energy  $E$ ,  $V(r)$  will only come into consideration as a small



“disturbance”, or “perturbation”; if we neglect it, we have as a solution of

$$\left\{ \frac{\hbar^2}{8\pi^2m} \Delta + E \right\} \psi = 0$$

the plane wave  $\psi = e^{(2\pi i/\hbar)pz}$ , where

$$E = \frac{p^2}{2m},$$

and the direction of the wave normal is assumed arbitrarily as parallel to the  $z$ -axis. The disturbance can be taken into account to a first approximation by substituting the plane wave expression for  $\psi$  in the term  $V(r)\psi$  in the original equation; we have then to find a solution (Born, 1926) of

$$\left( \frac{\hbar^2}{8\pi^2m} \Delta + E \right) \psi = V(r)e^{(2\pi i/\hbar)pz},$$

corresponding to a wave receding from the nucleus. It is perfectly clear, especially in the light of the analogy with the scattering of light waves, that the intensity of the secondary wave gives the number of electrons, belonging to a given incident beam, which are deflected in a definite direction; and this in effect implies the statistical interpretation as stated. A more rigorous investigation will be found in Appendix XX (p. 304); it is there shown how the intensity, i.e. the number of particles in a stream, per square centimetre and per second, is to be defined. In particular, if  $V(r)$  is chosen so as to correspond to a (screened) Coulomb field, the result obtained is precisely Rutherford's law of scattering (p. 53). As a matter of fact it is only for swift particles that this proof is valid; but it can be shown (Gordon, Mott, 1928) that the result is strictly correct. The exact solution differs from the approximate one only in terms which have no influence on the stream-intensity. This is very remarkable, and is analogous to the fact that, in the Coulomb field, the discrete term values, as given by wave mechanics, are in agreement with the values calculated with the help of quantised classical orbits.

If we extend this statistical interpretation to the case of discontinuous states, and if  $E_n$  is the energy and  $\psi_n$  the proper function of such a state, then  $|\psi_n|^2 dv$  is the probability that an electron will be found precisely in the volume element  $dv$ ; this holds in spite of the fact that the experiment if carried out would destroy the connexion with the atom altogether. The probability of finding the

electron somewhere or other in the atom must according to this interpretation be equal to 1; that is to say, the factor in the solution of the (homogeneous) wave equation, which in the first instance is quite indefinite, must be determined so that the equation

$$\int |\psi_n|^2 dv = 1$$

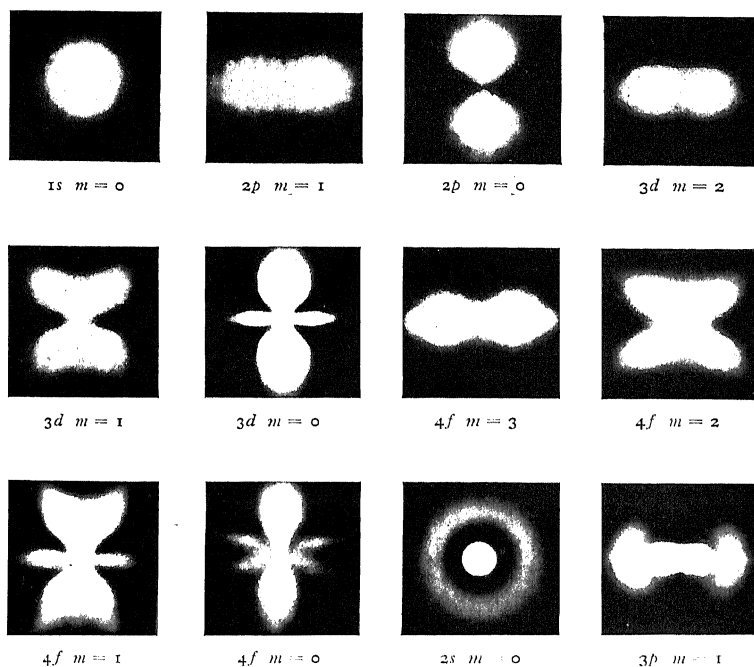
is satisfied. This "normalizing integral", which has no meaning except in the domain of discontinuous energy values, plays a remarkable part in this respect, that it does not vary with the time, even when we do not confine ourselves to stationary states alone, but substitute for  $\psi$  any solution at all of the wave equation, in the form containing the time.

We speak frequently of a *density distribution* of the electrons in the atom, or of an *electronic cloud* round the nucleus. By this we mean the distribution of charge which is obtained when we multiply the probability function  $|\psi_n|^2$  for a definite state by the charge  $e$  of the electron. From the standpoint of the statistical interpretation its meaning is clear; it can be represented pictorially in the way shown in fig. 21, Plate VII. The figures represent the projections (shadows) of the electronic clouds in various states; the positions of the nodal surfaces can be recognized in them at once.

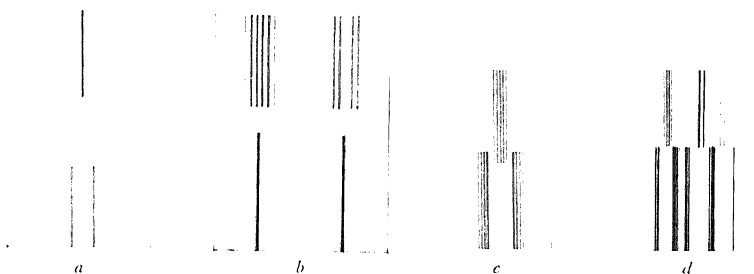
From another point of view, the statistical interpretation of wave functions suggests how the radiation emitted by the atom may be calculated on wave-mechanical principles. In the classical theory this radiation is determined by the electric dipole moment  $\mathbf{p}$  of the atom, or rather by its time-rate of variation. By the correspondence principle, this connexion must continue to subsist in the wave mechanics. Now the dipole moment  $\mathbf{p}$  is easily calculated by wave mechanics; if we adhere to the analogy with classical atomic mechanics, it is given by

$$\mathbf{p} = e \int \mathbf{r} |\psi_n|^2 dv = e \int \mathbf{r} \psi_n^* \psi_n dv,$$

where  $\mathbf{r}$  stands for the radius vector from the nucleus to the point of integration, or field point. (As usual, the asterisk denotes replacement by the conjugate complex quantity.) The integral represents of course the position of the "electrical centroid of the electronic cloud". Now, as is easily proved, this integral vanishes for all states of an atom so that the derivative of the dipole moment vanishes, and accordingly the emitted radiation also; that is, a stationary state does not radiate. This gives an explanation of the fact—unintelligible from the standpoint of Bohr's theory—that an electron which is revolving about the



Ch. V, Fig. 21.—The figures show graphically in the form of silhouettes the distribution of charge which, according to Schrödinger's theory, must exist in the electronic cloud. The nodal surfaces corresponding to the quantum numbers stated can be clearly seen, especially in the  $s$ -states, and in the states for which  $m$  is small.



Ch. VI, Fig. 2. Splitting patterns in a magnetic field. (a) Normal Lorentz triplet for the Cd line 6430 (the vibrations in the lower lines are perpendicular to the magnetic lines of force, those in the upper lines parallel to them). (b) Anomalous Zeeman effect for the Na  $D$ -lines (below without, above with magnetic field). (c) and (d) show Zeeman effects for Cr lines (above, vibrations parallel to field; below, perpendicular to it).



nucleus, and according to the classical laws ought to emit radiation of the same frequency as the revolution, can continue to revolve in its orbit without radiating. In wave mechanics this absence of emitted radiation is brought about by the fact that the elements of radiation, emitted on the classical theory by the individual moving elements of the electronic cloud, annul each other by interference.

But now, in analogy with the probability function or density function  $\psi_n^* \psi_n$  of a definite state, as defined above (p. 131), we can form, in the first instance in a purely formal way, the "transition density"  $\psi_m^* \psi_n$  corresponding to a transition from a state  $n$  to another state  $m$ ; it corresponds physically to the well-known "beat phenomena", which occur when two vibrations with neighbouring frequencies are superimposed on each other; its rhythm is given by the time factor

$$e^{-(2\pi i/\hbar)(E_n - E_m)t}$$

of the transition density, and the beat frequency is found from the difference of the energies of the two states:

$$\nu_{nm} = \frac{E_n - E_m}{\hbar}.$$

We now form also, in an analogous way to this, the dipole moment corresponding to the transition from  $n$  to  $m$

$$\mathcal{P}_{nm} = e \int \mathbf{r} \psi_m^* \psi_n dv = e \mathbf{r}_{nm} e^{-2\pi i \nu_{nm} t}.$$

It oscillates with the beat frequency given above, and therefore, by the classical formulæ (Appendix VIII, p. 275), radiates per unit time the energy

$$J = \frac{2}{3c^3} \ddot{\mathcal{P}}_{nm}^* \ddot{\mathcal{P}}_{nm} = \frac{2}{3} \frac{e^2}{c^3} (2\pi \nu_{nm})^4 |\mathbf{r}_{nm}|^2.$$

The quantity  $\mathbf{r}_{nm}$ , whose meaning follows from the formula for  $\mathcal{P}_{nm}$  (splitting of the time factor), is called the *matrix element of the vector co-ordinate  $\mathbf{r}$* ; as Schrödinger has shown, it is identical with the matrix element which, in Heisenberg's co-ordinate matrix, occupies the  $n$ th row and the  $m$ th column (Appendix XV, p. 291).

Thus, according to wave mechanics, we obtain the emitted radiation by calculating, in purely correspondence fashion, the radiation emitted by an oscillating dipole, as obtained by the rules of classical electrodynamics. From this it follows automatically that in the spectrum only those lines can occur whose frequency agrees with a beat

frequency between two states of the atom. These are exactly the lines to explain which Bohr had to introduce, as a fundamental postulate of his theory, the radiation condition

$$\nu_{nm} = \frac{1}{h} (E_n - E_m),$$

which is perfectly unintelligible from the classical standpoint. It is not to be understood, however, that both states  $n$  and  $m$  are excited simultaneously when emission occurs; it is rather a matter of their virtual presence. In point of fact, in order that a spectral line may be spontaneously emitted, the upper state must be excited in some way or other; the emission is then a companion process to the jump to the lower quantum state (vibration of an associated virtual resonator).

The intensity of the spectral line is the product of two factors, the number of excited atoms and the radiating strength  $J$  of an individual atom, which we have just calculated. Thus, with regard to the conditions of excitation of lines, those ideas in Bohr's theory which are brilliantly verified by experiment are just the ideas which are retained in their entirety in the wave mechanics. The latter theory adds a more exact calculation of the intensity  $J$  of the individual elementary act, depending on evaluation of the integrals occurring in the matrix elements, while on this question Bohr's theory could only with difficulty make a few statements, with the help of very considerable use of the correspondence principle.

As is shown in Appendix XXI (p. 308), the evaluation of the matrix elements in the case of the hydrogen atom leads to the selection rules  $\Delta l = \pm 1$  and  $\Delta m = 0, \pm 1$ ; that is, all matrix elements vanish which do not correspond to one of the transitions mentioned, and with them vanishes also the radiation of the corresponding frequencies. Application of this to hydrogen-like atoms (such as the halogens) gives the theoretical foundation for the facts that, for instance, transitions occur between  $s$  and  $p$  terms or between  $p$  and  $d$  terms, but not between  $s$  and  $d$  terms or between  $p$  and  $f$  terms.

Besides the discontinuous states there are also states forming a continuous range (with positive energy); they correspond to the hyperbolic orbits of Bohr's theory. The jumps from one hyperbola to another or to a stationary state give rise to the emission of the continuous X-ray spectrum emitted when electrons are scattered or caught by nuclei. The intensity of this spectrum has been calculated by Kramers (1923) from the standpoint of Bohr's theory by a very ingenious application of the correspondence principle. His

results can now be confirmed and improved by evaluating the matrix elements belonging to two states, one or both of which are in the continuous range (Oppenheimer, 1926).

Further, the method just explained for the wave-mechanical calculation of the radiation emitted can be put on a rigorous basis, as has been shown by Dirac. For this purpose we must on the one hand "quantise" the radiation field, and on the other hand take fully into account the coupling, as given by electrodynamics, between the radiation field and the atoms present in it. It is beyond our scope, however, to take up these problems here.

It might appear that wave mechanics involves a one-sided preference for the wave standpoint, and that the introduction of the corpuscle concept is therefore only made possible by the artificial importation of the statistical interpretation. As against that, it may be remarked, while the matrix or "quantum mechanics" of Heisenberg, briefly outlined above, is in complete agreement with wave mechanics in its content, and only differs from it in the form of presentation, still in its methods it attaches itself rather to corpuscular mechanics.

The true philosophical import of the statistical interpretation has already been explained in § 7 (p. 82). It consists in the recognition that the wave picture and the corpuscle picture are not mutually exclusive, but are two complementary ways of considering the same process—a process whose accessibility to intuitive apprehension is never complete, but always subject to certain limitations given by the principle of uncertainty. Here we have only one more important point to mention. The uncertainty relations, which we have obtained simply by contrasting with one another the descriptions of a process in the language of waves and in that of corpuscles, may also be rigorously deduced from the formalism of quantum mechanics—as exact inequalities, indeed; for instance, between the co-ordinate  $q$  and momentum  $p$  we have the relation

$$\Delta q \Delta p \geq \frac{h}{4\pi},$$

if  $\Delta q$  and  $\Delta p$  are defined as root-mean-squares (see Appendix XXII, p. 312).

## CHAPTER VI

# Spin of the Electron and Pauli's Principle

### 1. Alkali Doublets and the Spinning Electron.

The great success of Bohr's theory and especially of wave mechanics shows that in the interpretation of atomic processes we are on the right road to knowledge. The theory, however, as we have repeatedly emphasized in the preceding chapter, is still incomplete. In particular, explanations cannot yet be given of the anomalous Zeeman effect, the synthesis of the shells in the atom, &c. The wave mechanics of the atom, in its form as developed up to this point, still needs a far-reaching extension by the introduction of new ideas and hypotheses, which will form the subject of the present chapter.

The starting-point is given by the observation that the lines of the principal series in the alkalis are double. A well-known example is the *D*-line of sodium, whose doublet nature can be observed even with simple spectroscopic appliances. The splitting of the line is rather considerable—it amounts to 6 Å.; the two components are denoted by  $D_1$  and  $D_2$ , their wave-lengths being  $\lambda = 5896$  Å. for  $D_1$ , and  $\lambda = 5890$  Å. for  $D_2$ . The term analysis of the alkali spectra, to which the spectrum of sodium belongs, gives the definite information about them that the

s-                      terms ( $l = 0$ )                      are simple,  
p-, d-, . . . terms ( $l = 1, 2, \dots$ ) are double.

This experimental result cannot be explained either from the standpoint of Bohr's theory, or from that of wave mechanics so far as developed above. We have investigated above (p. 125) the most general motion of an electron in the atom, on the basis of its three degrees of freedom, and have arrived at the conclusion that the motion is completely determined and described by the three quantum numbers  $n, l, m$ . Any further splitting of the energy terms than that conditioned by these quantum numbers is therefore unintelligible, so



long as we adhere to the idea that the motion of the electron is at most triply periodic.

Under the compulsion of experiment, Uhlenbeck and Goudsmit (1925) put forward the following bold hypothesis. Previously the electron was regarded as an unextended structure, possessing only three degrees of freedom and, moreover, determined by two constants, viz. its charge  $e$  and its mass  $m$ . If, however, the electron is regarded as a structure with finite extension, then, like every extended system, it possesses three rotational degrees of freedom besides its three translational freedoms. It is therefore an obvious suggestion to ascribe to the electron angular momentum about an arbitrary axis (also frequently called "mechanical moment"). Developing the idea further, we must assume that the electron also possesses magnetic moment; for the electron carries an electric charge  $e$ , which rotates convectively, about the axis mentioned, along with the electron; so that the rotating electron is equivalent magnetically to a system of circular currents round the axis of rotation, and these, as we know, give rise to a magnetic moment. With regard to the magnitude of the magnetic and the mechanical moments, experiment must of course decide in the first instance; afterwards, we can try to deduce these magnitudes theoretically.

This property of the electron, in virtue of which it has a mechanical and a magnetic moment, is called its *spin*.

The magnitude of the *mechanical moment* follows immediately from known facts about the spectra of the alkalis. The angular momentum of the electron must of course, like every angular momentum, be quantised, and the same holds good for its component in a specially distinguished direction (external magnetic field). If then the value of the mechanical spin-moment is  $s$  (in units  $\hbar/2\pi$ ), there must, by the rules for the quantisation of direction, be  $2s + 1$  possible "settings" (i.e. orientations or inclinations) with respect to the special direction; the individual components of  $s$ , which we call  $\sigma$ , differing from each other by successive units. To see this, consider the analogous relations in the Bohr atom, in which the plane of an orbit with angular momentum  $l$  has precisely  $2l + 1$  possible settings with respect to the special direction, these settings being characterized by the components  $m$  of  $l$  in that direction (see fig. 10, p. 109). This extension of the concept of orbital angular momentum to spin angular momentum is justified by the fact that the consequences deduced from it are found to be in full accord with the facts. As regards this point, we may refer here once again to what was said in § 5, p. 127,

about the applicability of the classical vector-model to the description of atomic states in wave mechanics. In this chapter we shall be concerned almost exclusively with the conceptual scheme, and shall therefore use the pictorial vector-model, postponing the wave-mechanical treatment of the spinning electron until the close of the chapter (§ 8, p. 169).

As has just been brought out, the spin-moment  $s$  of the electron must have  $2s + 1$  possible settings with respect to a specially distinguished direction. Now experiment shows that the terms of sodium, excepting the  $s$ -terms, are double. This

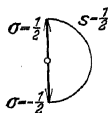


Fig. 1.—Setting of the spin with respect to a specially distinguished direction; there are two possible settings, parallel and antiparallel to this direction.

compels the conclusion that the spin-moment has only two possible settings (fig. 1), unless we are altogether wrong in assuming that this term-splitting is determined by the spin. Hence we must have  $2s + 1 = 2$ , or

$$s = \frac{1}{2}$$

(in units  $\hbar/2\pi$ ). The two possible settings have then the components

$$\sigma_1 = +\frac{1}{2}, \quad \sigma_2 = -\frac{1}{2}.$$

The occurrence of half-integers here as “quantum numbers” contradicts, at first sight, our ideas regarding the quantisation of angular momenta. It is to be noted, however, that the hypothesis we have used above of a rotating electron, extended in space, possesses merely heuristic value; we must be prepared, on following out these ideas, to encounter difficulties. (For instance, a point at the surface of the electron would have to move with a velocity greater than that of light, if such values as have been determined experimentally for angular momentum and magnetic moment are to agree with those calculated by the classical theory.) The use of half-integral components of angular momentum for the spin consistently leads, however, to results which are in complete agreement with the experimental facts. On the other hand, the wave mechanics of the spinning electron, in the form given to it by Dirac, leads automatically to this half-integral property, merely as a consequence of the conditions of linearity and relativistic invariance, without any subsidiary assumption.

An electron revolving about the nucleus possesses an orbital moment  $l$ ; besides this, it has the mechanical spin-moment  $s$ . The question arises: how are these two moments to be compounded with

each other? Bohr's theory would reply that they must be combined by the method of vector addition. This same rule for the composition of  $l$  and  $s$  holds good according to wave mechanics, although the proof (Wigner, v. Neumann, 1927) requires advanced mathematical methods (group theory). They therefore combine vectorially, giving a resultant (or total) moment  $j$  (in units  $\hbar/2\pi$ ), so that

$$\vec{j} = \vec{l} + \vec{s}.$$

After Sommerfeld,  $j$  is called the "*inner*" quantum number; it represents the total mechanical moment of the atom. It too must of course be quantised. Since  $s = \frac{1}{2}$ , there are only the two possibilities,

$$j_1 = l + \frac{1}{2}, \quad j_2 = l - \frac{1}{2};$$

$j$  is therefore half-integral in this case. For each  $l$ -value there are accordingly two possible values for the total mechanical moment, so that the corresponding terms are double. The  $s$ -terms alone ( $l = 0$ ) form an exception; they are always single, for in this case the only allowable value is  $j = s = \frac{1}{2}$ , since  $j$ , the total angular momentum, must always be positive. The double possibility for the setting of the spin with respect to the orbit is equivalent to a splitting of the energy terms, on account of the magnetic coupling of spin and orbit. The magnitude of the splitting is in fact given directly by the energy which is needed in order to turn the spin round, from one setting relative to  $l$  in the magnetic field of the orbit, into the other setting.

We take as an example the case with which we began (p. 136), that of the sodium  $D$ -lines. The term analysis shows that the upper state is a  $p$ -term, while the lower is an  $s$ -term. The former is double, corresponding to the two possible values of the total angular momentum  $j = \frac{1}{2}$  and  $j = \frac{3}{2}$ ; the lower term, being an  $s$ -term, is single ( $j = \frac{1}{2}$ ). The  $D_1$  line corresponds to the transition from the  $p$ -term with the inner quantum number  $j = \frac{1}{2}$ , the  $D_2$  line to the transition from the term with  $j = \frac{3}{2}$ .

The rule of vector addition can also be applied to the case of several electrons; in this case the orbital moments  $l_1, l_2, \dots$  of the individual electrons, and their spin-moments  $s_1, s_2, \dots$ , are compounded so as to give the total angular momentum  $j$ . Here  $j$  is integral or half-integral according as the number of electrons is even or odd. Similarly, the projection  $m$  of  $j$  on a specially distinguished direction can also be either integral or half-integral.

In conclusion, we may also recall the fact that the whole vector

figure of mechanical moments rotates round the direction of  $j$ , with uniform angular velocity; as follows from the meaning of  $j$ , the total angular momentum. This implies, as we have fully explained in the preceding chapter (p. 101), that the emitted radiation is subject to the rules  $\Delta j = \pm 1$ . In addition, however, as the theory agrees with experiment in showing, there are transitions with  $\Delta j = 0$ ; these correspond to changes of state in which the total angular momentum does not change. The fact that these transitions are permitted, while those with  $\Delta l = 0$  (or  $\Delta k = 0$ , see foot of p. 111) are forbidden, is capable of explanation on correspondence principles. We shall not, however, consider the matter more closely.

## 2. The Anomalous Zeeman Effect.

It has been shown in the last section, and will further appear in succeeding sections, that the hypothesis of the "spinning electron" has made it possible to understand the splitting of terms (multiplets), a phenomenon which the orbital picture by itself was quite incapable of explaining. The phenomenon in fact depends upon the possession of angular momentum by the electron itself; this internal angular momentum, by the quantum rules, can be directed in different ways with respect to the direction of the orbital moment, or with respect to a direction marked out by external means.

We shall now show that the electron's own magnetic moment, which is bound up with its mechanical moment, supplies the explanation of the *anomalous Zeeman effect*, i.e. the observed phenomenon that in a (weak) magnetic field a spectral line is split up into a considerable number of lines (fig. 2, Plate VII); while, according to classical theory, and also according to wave mechanics when spin is not taken into account, we can only have the *normal Zeeman effect*, i.e. the splitting up of every spectral line into a Lorentz triplet.

We may briefly recall the explanation of the normal Zeeman effect. The revolution of the electron produces a *mechanical moment*  $p_l$  of the orbital motion, and this is quantised by known rules:

$$p_l = l \frac{h}{2\pi}.$$

On the other hand, the revolving electron acts like a circular current of strength  $I = e(\omega/2\pi)$ , where  $\omega$  is the frequency of revolution, and so generates a magnetic field. But the magnetic field of a circular current  $I$  is, as we know, equivalent to that of a magnetic dipole of moment  $M = AI/c$ , where  $A$  is the area enclosed by the circuit, and

$c$  is the velocity of light. Hence the revolving electron behaves magnetically like a magnetic dipole of moment  $\pi r^2(e/c)\omega/2\pi$ ; since, however, the orbital angular momentum is  $p_l = \mu r^2\omega = l(\hbar/2\pi)$ , the magnetic moment  $M_l$  of the orbital motion becomes

$$M_l = \frac{eh}{4\pi\mu c} l = \frac{e}{2\mu c} p_l.$$

The value  $eh/(4\pi\mu c)$  therefore represents the smallest unit for the magnitude of a magnetic orbital moment in the atom; it is called the *Bohr magneton*.

If a homogeneous magnetic field is applied, the atom is set into precessional motion (fig. 3) about the direction of the field, as has been explained above (p. 109); consequently the component  $m$  of  $l$  in this direction must be a whole number (quantisation of direction). As for the supplementary

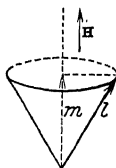


Fig. 3.—Precession of the orbital angular momentum round the direction of the magnetic field (in the absence of the spin it would always lead to the normal Zeeman effect).

energy, by which the energy of the atom is increased owing to the magnetic field, this is given by

$$E_{\text{magn}} = -M_l H \cos \theta,$$

where  $\theta$  is the angle between the magnetic field and the direction of the magnetic moment, i.e. the direction of  $l$ . But  $\cos \theta$  is obviously equal to  $m/l$ , so that

$$E_{\text{magn}} = -\frac{eh}{4\pi\mu c} H m.$$

The terms therefore split up in the magnetic field, the separation being

$$h\nu_{\text{magn}} = \frac{eh}{4\pi\mu c} H = h\nu_L,$$

where

$$\nu_L = \frac{e}{4\pi\mu c} H = 1.40 \times 10^6 H \text{ sec.}^{-1},$$

which is the same as the Larmor frequency already introduced (p. 109), i.e. the amount by which the frequency of a vibrating electrical system is changed in a magnetic field, according to the classical theory.

Thus every term splits up into  $2l + 1$  equidistant terms, corresponding to the  $2l + 1$  setting possibilities. In the emitted radiation,

however, every line should be split up into only three components, since the precessional motion is purely periodic, and therefore the selection rules  $\Delta m = 0, \pm 1$  come into play. In this way, therefore, we get only the *normal Zeeman effect* (see fig. 11, p. 110).

Even taking the spin of the electron into account, nothing in these relations would be changed, if we associated with the electron a magnetic moment, bearing to the mechanical spin-moment  $p_s = s(\hbar/2\pi)$  the same ratio as the magnetic orbital moment does to the mechanical, i.e. if

$$\frac{M_s}{p_s} = \frac{M_l}{p_l} = \frac{e}{2\mu c}.$$

For the total angular momentum would then be  $j$ , and the total magnetic moment would be  $M_j = (e\hbar/4\pi\mu c)j$ ; thus  $j$  and  $M_j$  would have the same direction, and would set themselves in the magnetic field in accordance with the quantisation of direction, or precess round the field direction in common. The single difference, with spin, would be that now, not  $2l + 1$ , but  $2j + 1$  setting possibilities exist, and that therefore every undisturbed term is split up by the magnetic field into  $2j + 1$  terms, but in such a way that the amount of the splitting would be exactly the same as before; in the spectrum there would be no difference at all.

The *anomalous Zeeman effect*, however, can be explained completely by assuming that the magnetic spin-moment is got from the mechanical, not by multiplying by  $e/2\mu c$ , as with orbital moments, but by multiplying by  $e/\mu c$ , so that

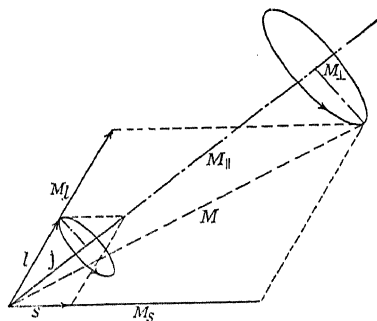
$$M_s = 2 \frac{e\hbar}{4\pi\mu c} s.$$

Since the mechanical spin-moment is always  $s = \frac{1}{2}$ , it follows that the magnetic moment of the electron is exactly equal to a Bohr magneton  $e\hbar/4\pi\mu c$ . This difference in the behaviour of the spin-moments as compared with the orbital moments can be put on a theoretical basis, as has been first shown by Thomas (1926), and recently in a much simpler way by Kramers (1935); it is a necessary consequence of the theory of relativity. Moreover, the connexion between mechanical and spin-moment, in the form we are now considering, follows rigorously from the relativistic equation of Dirac.

It is just this difference between the orbital and the spin-moments which is responsible for the anomalous Zeeman effect. The result of it is that the vector sum of the magnetic moments, i.e. the total magnetic moment  $M$ , is not in general in the same direction as the total

mechanical moment  $j$ . In fig. 4 this is shown for the case of one electron, but the principle is the same for the case of several electrons also. For clearness, the magnetic orbital moment  $M_l$  is shown twice as large as the mechanical orbital moment  $l$ ; hence, by the preceding, the magnetic spin-moment  $M_s$  must

Fig. 4.—Vector model for the anomalous Zeeman effect. The direction of the total angular momentum does not coincide with the direction of the resultant magnetic moment; only the component  $M_{\parallel}$  parallel to  $j$  is magnetically effective; the other component  $M_{\perp}$  disappears when averaged, on account of the rotation of the vector figure about  $j$  (the total angular momentum).



be shown four times as large as the mechanical spin-moment  $s$ . The resultants  $M$  and  $j$  therefore fall in different directions.

In accordance with the meaning of  $j$ , the total angular momentum, we must regard the atom, and with it the whole vector figure, as in rotation about the direction given by  $j$ ; any vector, not in this direction, therefore precesses round it. On account of the high frequency of the precessional motion (it can be shown that the  $h\nu$  corresponding to this frequency is of the order of magnitude of the "fine structure splitting" of the terms, which is determined by the coupling of  $l$  and  $s$  to the vector  $j$ ; or, in the case, say, of the sodium  $D$ -lines, of the order of magnitude of the splitting of the  $p$ -terms, i.e. approximately,  $\nu \sim 5 \times 10^{11}$  cm.<sup>-1</sup>) only time-averages can come into question, for magnitudes which vary slowly as compared with this frequency. Thus, for example, the atom will behave magnetically, in the presence of an external field, as if the atom had the magnetic moment  $M$ , where the bar indicates a time-average. But the time-average of  $M$  is equal to the projection of  $M$  on the axis of rotation, i.e. equal to the component  $M_{\parallel}$ ; the component  $M_{\perp}$  perpendicular to the axis disappears on formation of the time-average.

In the presence of a (weak) external field the atom therefore possesses an effective magnetic moment  $M_{\parallel}$  in the direction of  $j$ . On account of the angular momentum it precesses about the direction of the field, and the same considerations now hold as those adduced above:  $j$  possesses, in consequence of this precessional motion, only  $2j + 1$  possible settings with respect to the field direction, these being characterized by the component  $m$  of  $j$  in this direction. The magnetic energies of these settings are given individually by  $-M_{\parallel}Hm/j$ ;

the undisturbed term is therefore split up by the magnetic field into  $2j + 1$  terms with the separation  $M_{\parallel}H/j$ .

The difference now, as compared with the earlier case, is that formerly (in the normal Zeeman effect) the magnetic moment was equal for all terms to the Bohr magneton multiplied by the quantum number of the total angular momentum; the splitting was the same for all terms, irrespective of their quantum numbers, being equal to the Bohr magneton multiplied by the field strength  $H$ . When, however, the spin and the vector composition of the moments are taken into consideration, we obtain an "effective" magnetic moment of the atom, which in general is not given by the product of Bohr magneton and total angular momentum, but depends also on the rest of the quantum numbers and in particular on the angles occurring in the vector figure. If we write, purely formally in the first instance,

$$M_{\parallel} = \frac{eh}{4\pi\mu c} jg,$$

then the factor  $g$  gives the divergencies which occur in our vector model, as compared with the theory of the normal Zeeman effect. The additional magnetic energy is then given by

$$E_{\text{magn}} = -M_{\parallel}H \frac{m}{j} = -\frac{eh}{4\pi\mu c} Hmg = -h\nu_L mg,$$

where, as above,  $\nu_L$  is the classical Larmor frequency. The undisturbed term is therefore split up in the magnetic field, certainly again into

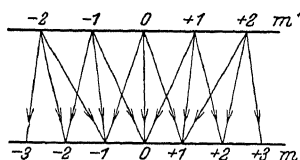


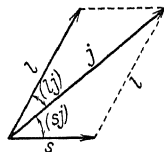
Fig. 5.—Transitions in the anomalous Zeeman effect; since the splitting is different in the various term groups, we get in general just as many separated lines as there are possible transitions altogether.

$2j + 1$  equidistant terms, but the amount of the splitting  $h\nu_L g$  is not equal to that in the normal Zeeman effect, viz.  $h\nu_L$ , but differs from it by the factor  $g$ , which is called, after its discoverer, the Landé splitting factor (1923). It varies from term to term, and is thus the determining cause of the anomalous behaviour of the atom in the Zeeman effect. The splitting pattern is not the Lorentz triplet; many more lines occur, in accordance with the fact that the energy differences corresponding to the selection rules  $\Delta m = -1, 0, 1$  are not now, as they are in the normal Zeeman effect, the same for all values of  $m$  (fig. 5).



The splitting up of a line in the anomalous Zeeman effect is therefore essentially determined by the Landé factors for the upper and lower states. These factors, as will now be shown, can be ascertained with comparative ease from the developments already given with regard to the vector model. For this purpose we have to calculate the value of

Fig. 6.—Vector composition of the orbital moment  $l$  and the spin-moment  $s$ , giving the total angular momentum  $j$ .



the component  $M_{\parallel}$  of the magnetic moment. From fig. 6 we can at once read off the relation

$$M_{\parallel} = M_l \cos(l, j) + M_s \cos(s, j).$$

If we now substitute for  $M_{\parallel}$  the expression used above,  $(eh/4\pi\mu c)jg$ , and for the magnetic moments of the orbit and the spin, the values already obtained, we find (note the factor 2 in the spin term!)

$$\frac{eh}{4\pi\mu c} jg = \frac{eh}{4\pi\mu c} \left\{ l \cos(l, j) + 2s \cos(s, j) \right\}$$

which gives

$$g = \frac{l}{j} \cos(l, j) + \frac{2s}{j} \cos(s, j).$$

The values of the cosines appearing in this equation may be written down at once from the triangle of the vectors  $l$ ,  $s$  and  $j$  (fig. 6):

$$\cos(l, j) = \frac{j^2 + l^2 - s^2}{2jl}; \quad \cos(s, j) = \frac{j^2 + s^2 - l^2}{2js}.$$

For the Landé splitting factor we thus obtain

$$\begin{aligned} g &= \frac{j^2 + l^2 - s^2}{2j^2} + 2 \frac{j^2 + s^2 - l^2}{2j^2} \\ &= 1 + \frac{j^2 + s^2 - l^2}{2j^2}. \end{aligned}$$

We have deduced this formula in accordance with the classical vector model representation. In quantum mechanics this representation is certainly still permissible, but with this difference, that the square of the magnitude of an angular momentum, with the quantum number  $l$ , is not equal to  $l^2$  as in Bohr's theory, but is given by  $l(l+1)$ . This is proved in Appendix XIX (p. 302) for the orbital angular momentum;

the same holds good for the spin angular momentum, so that we can adapt the formula for the splitting factor to wave mechanics, by writing everywhere  $j(j+1)$ ,  $l(l+1)$ ,  $s(s+1)$ , instead of  $j^2$ ,  $l^2$ ,  $s^2$ :

$$g = 1 + \frac{j(j+1) + s(s+1) - l(l+1)}{2j(j+1)}$$

$$= \frac{3}{2} + \frac{s(s+1) - l(l+1)}{2j(j+1)}.$$

This method of deducing the  $g$ -formula may seem rather unconvincing, since classical ideas are applied throughout, and it is only at the very end, as it were, that corrections from wave mechanics are brought to bear. It may be mentioned, however, that the formula can also be obtained from wave mechanics directly, though only, it is true, with the assistance of group theory.

An immediate proof of the correctness of the formula is given, however, by experiment. In Appendix XXIII (p. 318) the splitting pattern in the anomalous Zeeman effect is calculated for the  $D$ -lines of sodium. Even in this simple example—an atom with one radiating electron—it can be seen how complicated, comparatively, is this splitting pattern in the anomalous Zeeman effect. Experiment, however, has fully confirmed the correctness of the calculated pattern, with complete spectroscopic accuracy. In every case in which the anomalous Zeeman effect has been investigated experimentally, theory and experiment have been found to be in complete agreement. Considering the enormous numerical material which has been collected up to the present time, and which the calculation of Landé's factors has invariably proved capable of explaining, one is compelled to conclude without going further that the theoretical interpretation of the anomalous Zeeman effect, and the results which it involves with respect to the magnitude of the mechanical and magnetic spin-moments, represent the actual truth.

In conclusion, we shall say something about one more phenomenon observed in the Zeeman effect; this is called, after its discoverers, the *Paschen-Back effect* (1921). We have stipulated above that the magnetic field is not to be too strong. This stipulation was tacitly used when we replaced the magnetic moment by its time-average  $\bar{M} = M_{||}$ . This is justified so long as the rotation of the whole vector model about the direction of the angular momentum  $j$  is much more rapid than its precessional motion about the direction of the magnetic field, the frequency of which is, approximately,  $\nu_L = 1.40 \times 10^6 H \text{ sec.}^{-1}$ .

For the former, in the case of the splitting of the sodium terms, we found (p. 136)  $\Delta\lambda \sim 6 \text{ \AA.}$ , corresponding to  $\Delta\nu \sim 5 \times 10^{11} \text{ sec.}^{-1}$ , so that the required condition is certainly fulfilled in this example, for fields up to some thousand gauss.

If, however, we increase the magnetic field strength until the two frequencies are of the same order of magnitude, the foregoing considerations are no longer valid, since it is then not the time-average of  $M$ , but  $M$  itself, that is in question. We are now in the region of the Paschen-Back effect. We can describe it in another way by saying that the internal energy of precession about  $j$  becomes comparable with the external energy of precession of  $j$  about  $H$ , so that the fine structure splitting depending on the spin is of the same order of magnitude as the splitting of the terms in the magnetic field. If the magnetic field strength is raised further still, so that the energies depending on the setting in the field become much greater than the energy of coupling between the orbital moment  $l$  and the spin-moment  $s$ , we obtain the normal Zeeman effect; in fact, this coupling is then practically completely annulled, and the orbital and spin-moments

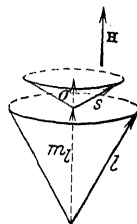


Fig. 7.—Vector model for the Paschen-Back effect (transition to the normal Zeeman effect). Since the energy of the orbital moment and spin-moment in the magnetic field is greater than the magnetic interaction between orbit and spin, the orbital and spin-moments precess separately round the field direction.

precess independently round the direction of  $H$  (fig. 7). The total magnetic energy is then given by

$$E_{\text{magn}} = -\frac{eh}{4\pi\mu c} H \left\{ l \cos(lH) + 2s \cos(sH) \right\}.$$

If we continue to denote by  $m_l$  (a whole number) the projection of  $l$  on the field direction, and observe that in this case the spin can only set itself parallel or antiparallel to the field ( $\sigma = \pm \frac{1}{2}$ ), it follows from the above equation and the half-integral character of the spin-moment that

$$E_{\text{magn}} = -\frac{eh}{4\pi\mu c} H(m_l \pm 1).$$

We therefore in this case obtain a term-splitting with the term-separation corresponding to the normal Zeeman effect. Thus when the magnetic field is steadily increased, a gradual transition takes place from the anomalous to the normal Zeeman effect; the transitional zone is referred to as that of the Paschen-Back effect. The nomen-

clature, it may be noted, is not appropriate—normally (at ordinary field strengths) we get the “anomalous” Zeeman effect, while the “normal” effect is only got at abnormally high field strengths.

### 3. The Hydrogen Atom and X-ray Terms.

We shall now investigate the question of how the statements of the atomic theory with regard to the values of the terms are to be corrected in the light of the recognition of the existence of spin. We begin with the simple example of the hydrogen atom, and with the schemes of terms which may be called hydrogen-like (alkali terms, X-ray terms). In Chapter V (pp. 112, 126) we discussed the values of

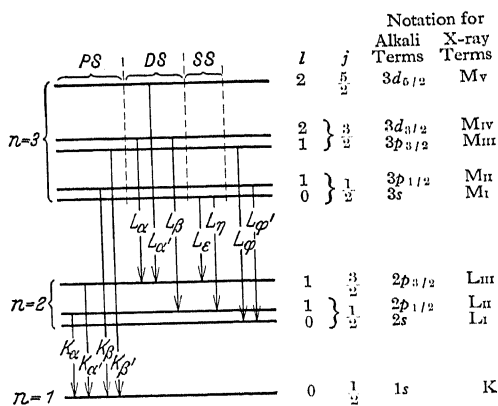


Fig. 8.—Diagrammatic synopsis of the notation for the various terms and lines in alkaline and X-ray spectra.

these terms fully, on the basis of Bohr's theory and wave mechanics, without taking into account the spin of the electron. We may state the results once more, in brief summary. The hydrogen-like spectra arise when an electron moves in a Coulomb field (hydrogen terms) or in a Coulomb-like central field (screening by the rest of the electrons; alkali and X-ray terms). The value of the term depends mainly on the principal quantum number: Balmer terms  $-R\hbar/n^2$ . If the field deviates from the Coulomb field a correction has to be applied, which depends not only on  $n$  but also on the azimuthal quantum number  $l$ ; it was denoted above (p. 107) by  $\epsilon(n, l)$ , where  $k = l + 1$ . A similar correction is required if we take into account the relativistic variability of mass, the influence of which on the hydrogen terms has been mentioned above (p. 107).

The existence of electronic spin leads to a further correction of

the terms, since it gives rise to an additional determining number  $j = l \pm \frac{1}{2}$ . This additional number, as we have explained in § 1, p. 139, causes a splitting up of the terms into doublets; but the *s*-terms, and these alone, remain single, for in this case  $j$  can have only one value, viz.  $\frac{1}{2}$ . Instead of the scheme of terms in fig. 13, p. 112, we therefore obtain the scheme of fig. 8. The terms fall in the first place into widely separated groups corresponding to the principal quantum number  $n$  ( $n = 1, 2, 3, \dots$ ). The figure is not at all correct in its proportions; in reality the distance between two groups of terms, with different quantum numbers, is  $10^3$  to  $10^4$  times greater than the splitting within such a group; the transitions between two different groups of terms correspond in the optical spectra to wave-lengths of a few thousand Å., while the fine structure splitting of the lines amounts at most to a few Å.

Within the group of terms with the same principal quantum number, we have in a preceding section (p. 108) discriminated between the individual terms by specifying the azimuthal quantum number  $k$ ; instead of  $k$  we have now two quantum numbers, viz. the azimuthal quantum number  $l$  which corresponds to  $k - 1$ , and the inner quantum number  $j$ . Instead of  $l = 0, 1, 2, \dots$  (or, as formerly,  $k = 1, 2, 3, \dots$ ), it is customary, as was mentioned on p. 126, to use the letters *s, p, d, \dots*; the value of  $j$  is given as a suffix attached to these letters. Examples of this notation are shown in fig. 8, for the case of the alkali terms. For the X-ray terms a different notation has become established, the terms being distinguished by a capital letter specifying the shell in which they lie (*K* shell for  $n = 1$ , *L* shell for  $n = 2$ , and so on), and by a roman numeral attached to the letter as a suffix, and corresponding to the order of numbering within the shell.

The case of hydrogen is peculiar in one respect. Experiment gives distinctly fewer terms than are specified in the term scheme of fig. 8; for  $n = 2$  only two terms are found, for  $n = 3$  only three, and so on. The theoretical calculation shows that here (by a mathematical coincidence, so to speak) two terms sometimes coincide, the reason being that the relativity and spin corrections partly compensate each other. It is found that terms with the same inner quantum number  $j$  but different azimuthal quantum numbers  $l$  always strictly coincide, for instance, the *ns* and the  $np_{\frac{1}{2}}$  term, the  $p_{\frac{3}{2}}$  and the  $d_{\frac{3}{2}}$  term, and so on; such pairs of terms are drawn close together in fig. 8. For the value of the terms a formula was given by Sommerfeld (1916), even before the introduction of wave mechanics; the same formula is also obtained when the hydrogen atom is calculated by Dirac's relativistic

wave mechanics. This formula, which reproduces the values of the hydrogen terms with the greatest exactness, is:

$$E + E_0 = E_0 \left[ 1 + \frac{\alpha^2 Z^2}{(n_r + \sqrt{n_\phi^2 - \alpha^2 Z^2})^2} \right]^{-\frac{1}{2}}$$

Here  $E$  denotes the energy of the bound electron after deducting the rest energy, and  $E_0$  is the rest energy  $mc^2$ ;  $n_r$  is the radial quantum number;  $n_\phi$  (Sommerfeld) is identical with Bohr's azimuthal quantum number  $k$ , and corresponds therefore to the  $l + 1$  of wave mechanics. Since, however, as we have just seen, two terms with different  $l$  but the same  $j$  always coincide when we take the spin into account, discrimination between the terms by means of the quantum number  $n_\phi$  is identical with discrimination by means of  $j$ ; we therefore have  $n_\phi = j + \frac{1}{2}$ . The principal quantum number is then found as the sum  $n = n_r + n_\phi$ . The constant  $\alpha$  is given by

$$\alpha = \frac{2\pi e^2}{hc} \sim \frac{1}{137};$$

dimensional considerations show at once that it is a pure number, the only quantity indeed of zero dimensions (ignoring trivial numerical factors) which can be formed from the three atomic constants  $e$ ,  $h$  and  $c$ . Since it gives the amount of the fine structure splitting, it is called (after Sommerfeld) the *fine structure constant*.  $Z$  is the atomic number (1 for hydrogen, 2 for  $\text{He}^+$ , and so on). On account of the smallness of  $\alpha$ , Sommerfeld's formula can be expanded in ascending powers of  $\alpha^2 Z^2$ ; a simple calculation gives

$$E = -\frac{R h Z^2}{n^2} \left\{ 1 + \frac{\alpha^2 Z^2}{n^2} \left( \frac{n}{n_\phi} - \frac{3}{4} \right) + \dots \right\},$$

where  $R$  is written for  $E_0 \alpha^2 / 2h$ , so that  $R$  is simply the well-known Rydberg constant. The Balmer term  $-R h Z^2 / n^2$  is therefore modified by a correcting factor, which depends on  $n_\phi$  and gives the fine structure; the quantity  $\epsilon(n, k)$  of the formula on p. 107 is equal to the additive correction

$$-\frac{R h \alpha^2 Z^4}{n^4} \left( \frac{n}{n_\phi} - \frac{3}{4} \right).$$

It should be particularly emphasized that the X-ray terms also are well represented for all elements by Sommerfeld's formula.

One small remark on the fine structure constant may be added. When a physical formula is to be devised, it may be possible to go a

good length on the score of dimensional considerations alone. An example of this sort is the formula deduced in § 2, p. 46 for the scattering of light by electrons. Since the ratio of the total scattered energy  $J$  to the primary energy  $J_0$  incident per square centimetre has the dimensions of an area, and since from the three classical electronic quantities  $e$ ,  $m$ ,  $c$  (we are concerned here with a purely classical effect) only one quantity having the dimensions of a length can be formed, viz.  $e^2/mc^2$ , it follows that the scattering formula must be of the form  $J/J_0 = k(e^2/mc^2)^2$ , where  $k$  is a factor of proportionality of dimensions zero (and of order of magnitude 1). The real task of the theory is to determine the constant  $k$ , by means of geometrical considerations based upon definite ideas with regard to the mechanism of the process. The occurrence, in the theory of the splitting of terms, of the dimensionless fine structure constant  $\alpha$ , whose value is given by  $e$ ,  $c$  and  $\hbar$ , suggests the idea that there must be a deeper connexion between these three quantities, on the strength of which one of them can be determined in terms of the other two. If this is so, there must be some process involved which we do not yet understand, its geometry and mechanism being still quite unknown. The discovery of this connexion is a problem which fascinates many physicists. Its solution would mean a great step towards understanding why charge, mass, energy, &c., occur in minimum elementary quantities. We may add that Eddington has put forward a theory for the geometrical explanation of  $\alpha$ , according to which  $1/\alpha$  is the value of  $\frac{1}{2}n^2(n^2+1)+1$  for  $n=4$ ; we cannot, however, go into this further. It may be remarked that the idea that  $1/\alpha$  is an integer is somewhat dubious, as the most exact experimental values give  $1/\alpha = 137.2$ .

We now proceed to consider the spectral lines which are possible, consistently with the above scheme of terms. The selection rules have already been stated (p. 140); they are

$$\Delta l = \pm 1, \quad \Delta j = 0, \quad \pm 1;$$

they do not allow all transitions, but only those indicated by arrows in the scheme of terms (fig. 8, p. 148).

In the hydrogen and alkali spectra, as we have already explained for the simple Bohr theory (fig. 13, p. 112; fig. 8, p. 148), the lines which correspond to transitions from a  $p$ -term to the ground state ( $s$ -term) are called the lines of the principal series; a principal series of higher order contains the lines which lead to the next  $s$ -term (in our diagram to the  $2s$ -term). Transitions from the  $d$ -terms to the  $p$ -terms give the diffuse series, those from the  $s$ -terms to the  $p$ -terms

the sharp series. The higher series (fundamental series, &c.) lead from higher levels to the terms of the third quantum state ( $n = 3$ ), and therefore fall outside the scope of our diagram of terms. We thus obtain—for the case of sodium—the diagram of terms and lines shown in fig. 9, which replaces the simple diagram of p. 112 (fig. 13).

For X-ray spectra an essentially different notation has secured general acceptance; this also is indicated in fig. 8, p. 148 (see also fig. 15, p. 113). The lines which correspond to transitions to the  $K$ -level are called  $K$ -lines; lines whose final state is a term of the  $L$  shell are called  $L$ -lines; and so on. Individual lines within these general classes are distinguished by small Greek letters in accordance with a recognized convention, which was also exemplified in fig. 15, p. 113. The  $K_\alpha$  and  $K_{\alpha'}$  lines form the doublet which arises in a transition of the electron from the  $L$  to the  $K$  shell;  $K_\beta$  and  $K_{\beta'}$  correspond to the transitions from the  $M$  to the  $K$  shell; and so on.

#### 4. The Helium Atom.

In the helium atom two electrons revolve about the nucleus (nuclear charge  $2e$ ); we have therefore 6 co-ordinates to deal with instead of 3, with the result that an exact solution is no longer possible. For the purpose of obtaining a general idea of the possible states, an exact solution is, however, not at all necessary; following Bohr, we can in the first place neglect the mutual interaction of the electrons, and for a first approximation treat the problem as if the two electrons moved undisturbed in the field of the nucleus. Afterwards, the interaction can be taken into account by the methods of the theory of perturbations.

We shall therefore (as with the hydrogen atom) associate three quantum numbers with each of the two electrons:

$$n_1, l_1, j_1; \quad n_2, l_2, j_2,$$

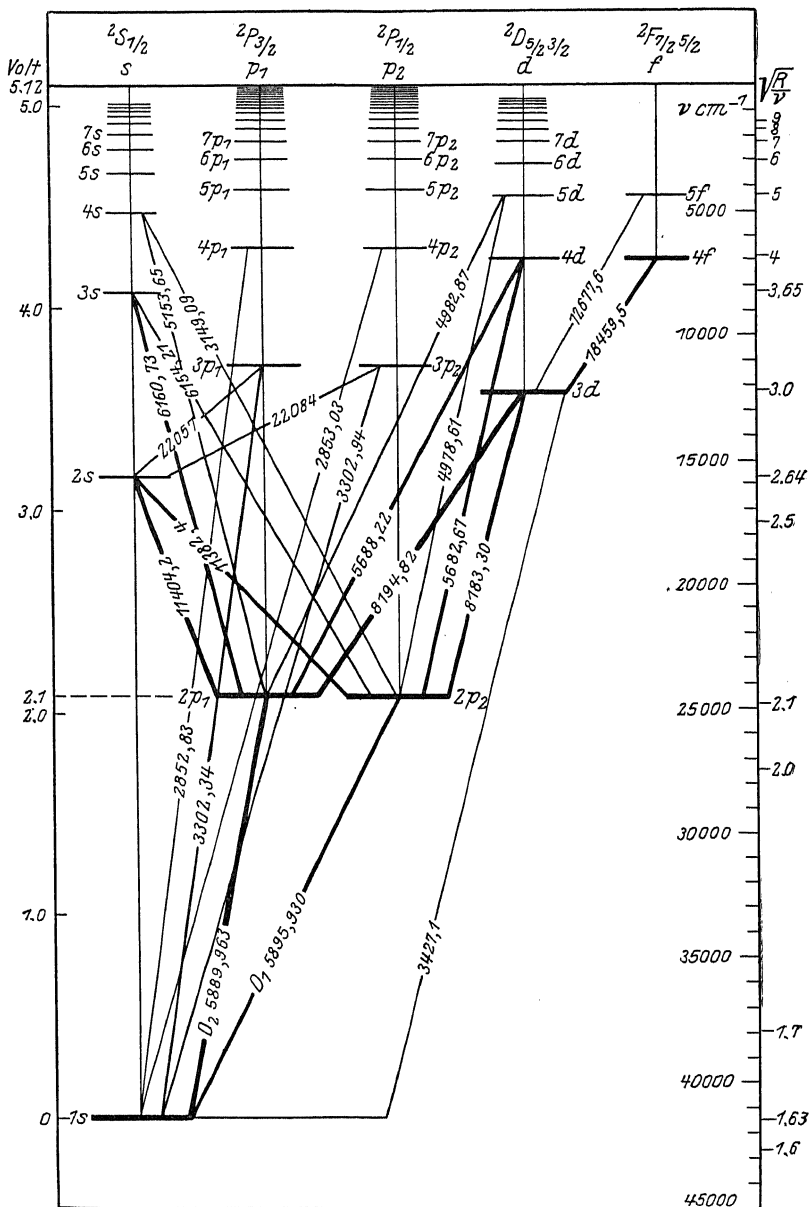
and we shall assume once for all that in the case of different  $l$ -values the first electron possesses the higher azimuthal quantum number ( $l_1 > l_2$ ). The corresponding angular momenta are compounded vectorially (neglecting the interaction between the two electrons); we thus obtain a total angular momentum

$$\vec{j} = \vec{j}_1 + \vec{j}_2,$$

as also total orbital and spin angular momenta given respectively by

$$\vec{l} = \vec{l}_1 + \vec{l}_2, \quad \vec{s} = \vec{s}_1 + \vec{s}_2.$$





If, however, the interaction of the two electrons is taken into account, the angular momenta of the electrons separately are no longer integrals of the equations of motion, so that the angular momentum vectors are not now fixed in space. The *total* angular momentum  $j$  is constant, however, and of course must still be quantised even for an arbitrary system of electrons, since it corresponds to a rotation of the electronic configuration as a whole. The question is now, whether a many-electron (two-electron) problem can still be characterized, at least approximately, by other angular momentum quantum numbers. The important factor here is the interaction of the two electrons, and the coupling relations between the individual angular momentum vectors (orbital and spin).

If in fact, in consequence of the interactions being slight, a precessional period is slow, it will be possible to associate the corresponding angular momenta, approximately, with quantum numbers, viz. those which the angular momenta would have, if there were no coupling. It is entirely a matter of the strength of the various couplings.

An obvious assumption to make

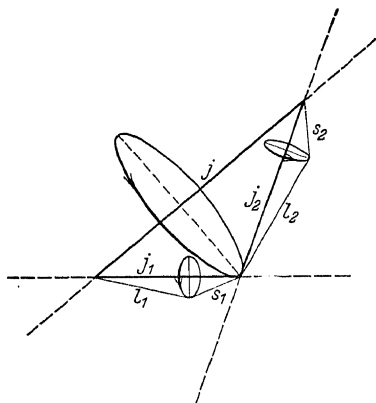


Fig. 10.—Scheme of the so-called (*jj*)-coupling: to a first approximation the orbital and spin-moments of each electron are compounded into a resultant; the vector sum of these two resultants gives the total angular momentum  $j$ .

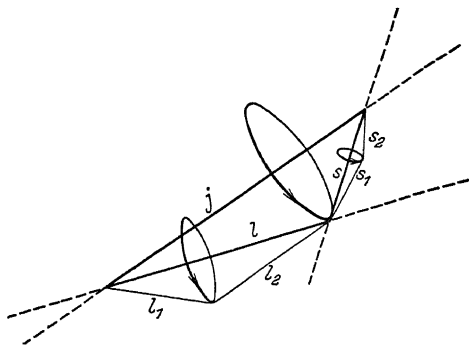
is the following. For every electron the orbital and spin moments are firmly coupled; but the various electrons influence each other comparatively little. Every electron will then be characterized individually by the quantum number of its angular momentum  $j$  (orbital plus spin); that is to say, the vector sum of  $l_v$  and  $s_v$ , for the  $v$ th electron, still approximately carries out a motion of precession round the direction of  $j_v$ , in spite of the disturbance by the other electrons; then the  $j$  vectors of the various electrons are compounded into a total angular momentum  $j$ , and in their turn carry out a precessional motion round the latter (see fig. 10).

Another limiting case of the coupling relations is that in which the spin vectors and the orbital vectors are compounded separately into resultants  $l$  and  $s$ , so that for two electrons we have the diagram of fig. 11. The vectors  $l_1$  and  $l_2$  rotate round the total orbital moment  $l$ ,

the spin vectors  $s_1$  and  $s_2$  round  $s$ , and then these two vectors precess about  $j$ . One can understand how this vector model may come into play, by supposing that besides the electrical forces of repulsion other forces of a magnetic nature also act between the electrons; the latter forces being determined by the magnetic moment of the electrons, and producing a strong coupling between the spin vectors of the two electrons.

Experimental results show that as a rule (in helium and the alkaline earths) the second case of coupling is the one which occurs, it is called the normal coupling case or, after its discoverers, the Russell-Saunders coupling. This is the only case we shall consider here although the case first described also actually occurs, as well as intermediate stages between these two extreme cases. We may remind the reader of the similar circumstances in the case of the abnormal Zeeman effect, where the

Fig. 11.—Scheme of the Russell-Saunders or (*ls*)-coupling; to a first approximation the orbital moments of the electrons are compounded into a total orbital moment  $l$ , also their spin-moments into a resultant spin-moment  $s$ , and finally  $l$  and  $s$  into the total angular momentum  $j$ .



strengthening of the coupling between the external magnetic field and the individual magnetic moments determines a transition to the normal Zeeman effect, via the Paschen-Back effect.

In helium, the coupling relations are normal. The two orbital moments  $l_1$  and  $l_2$  are therefore in the first place combined into a total orbital moment  $l$ , which must be a whole number and can therefore take only the values  $l = l_1 - l_2, l_1 - l_2 + 1, \dots, l_1 + l_2$ . Similarly, the two spin-moments  $s_1$  and  $s_2$  combine into the total spin-moment which (for two electrons) must likewise be a whole number: since  $s_1$  and  $s_2$  have the value  $\frac{1}{2}$ , the two values  $s = 0$  and  $s = 1$  are the only possible ones; in the former case the spins are antiparallel, in the latter parallel. (In fig. 11, for easier visualization we have assumed that the two spins form an obtuse angle with each other; in reality they must always be either parallel or antiparallel to each other.)

The term scheme of helium (see fig. 12) breaks up therefore into two separate groups, according as the total spin is zero or unity; in

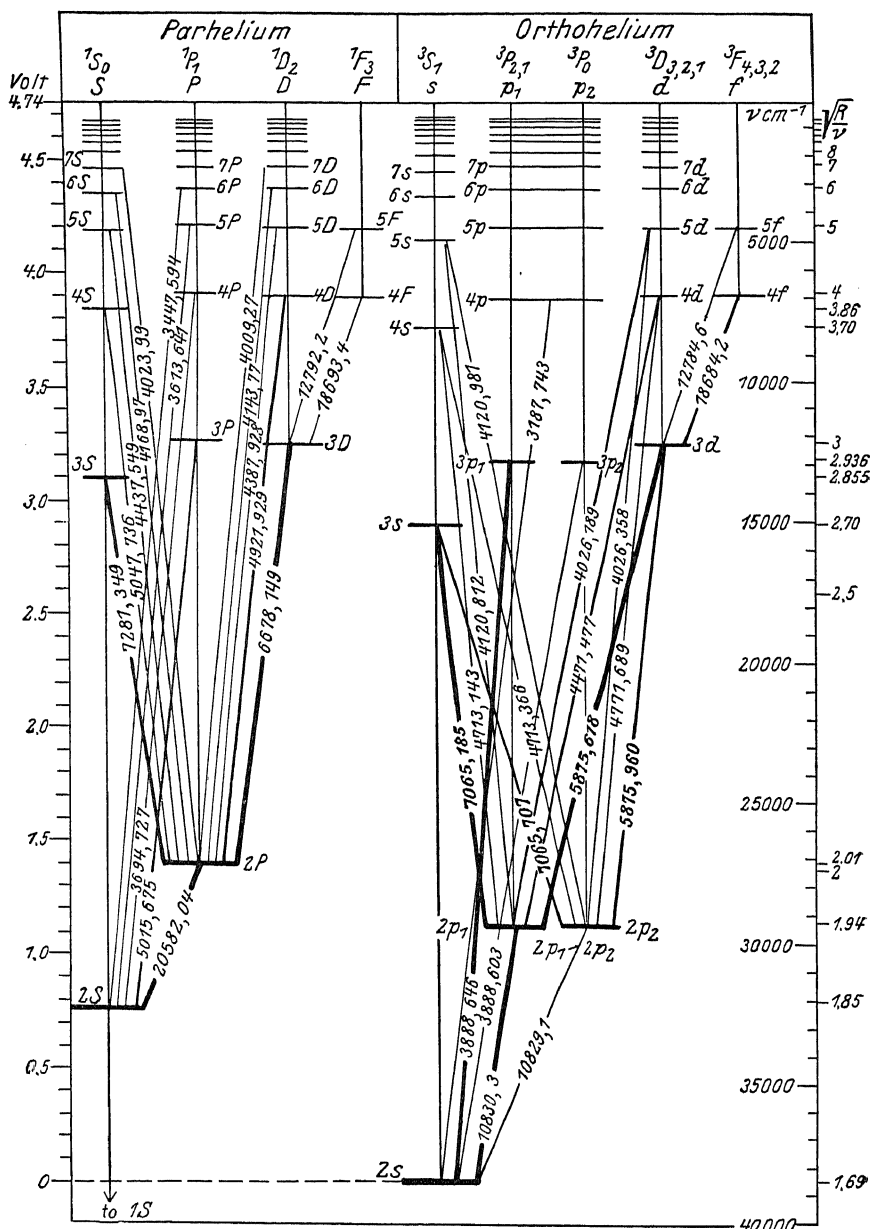


Fig. 12.—Term scheme for helium; there are no inter-combination lines between orthohelium and parhelium; orthohelium (for not too great dispersion) has the character of a doublet system. The  $1S$  term of parhelium is situated much deeper as indicated by the arrow

the former case we speak of *parhelium terms*, in the latter of *orthohelium terms*. This separation of the terms into two groups is justified by the fact, which has been proved experimentally and can easily be put on a theoretical footing, that in general there are no transitions between terms of the first group and terms of the second. From the physical standpoint it is accordingly permissible to distinguish sharply between parhelium and orthohelium, especially as it is only in exceptional cases that it is possible to transform a helium atom, which has once shown the parhelium spectrum, in such a way that it emits the lines of orthohelium.

In *parhelium*, then, the total spin moment  $s$  is 0; hence the total orbital moment is identical with the total angular momentum:  $j = l$ . This implies that the whole of the terms of parhelium are singlets, i.e. that to every azimuthal quantum number  $l$  there belongs only a single term with the inner quantum number  $j$  equal to  $l$ .

In *orthohelium*, on the other hand, the total spin moment is equal to 1, and it combines vectorially with the total orbital moment to form the total moment  $j$ . Since all three vectors are whole numbers,  $j$  in this case can take the three values

$$j = l - 1, \quad l, \quad l + 1.$$

There are therefore three terms for every azimuthal quantum number, and the term scheme of orthohelium is a triplet system. Here also, however, as in the doublet spectra discussed above (p. 136), the  $s$ -terms are single; in fact,  $l = 0$  for these, so that  $j$  can only be equal to  $s$ , i.e. to 1. In the alkaline earths (Be, Mg, Ca, Sr, Ba), which like helium have two external electrons, the triplet character can be easily observed in the spectrum. In helium, however, a special feature occurs: it possesses practically a doublet spectrum, since the  $P_1$  and  $P_2$  terms nearly coincide. As an example of this, we take the first line of the sharp series, which corresponds to the transition from  $3S$  to  $2P$ . Exact measurements by Houston have shown that this line consists of three components:

$$\left. \begin{array}{l} 3S \rightarrow 2P_0 \dots 7065.707 \text{ \AA}, \\ 3S \rightarrow 2P_1 \dots 7065.212 \text{ \AA}, \\ 3S \rightarrow 2P_2 \dots 7065.177 \text{ \AA}, \end{array} \right\} \Delta\lambda = 0.495 \text{ \AA},$$

$$\left. \begin{array}{l} 3S \rightarrow 2P_0 \dots 7065.707 \text{ \AA}, \\ 3S \rightarrow 2P_1 \dots 7065.212 \text{ \AA}, \\ 3S \rightarrow 2P_2 \dots 7065.177 \text{ \AA}, \end{array} \right\} \Delta\lambda = 0.035 \text{ \AA}.$$

We see from these numbers that the splitting between the two lower terms is only about one-fourteenth of that between either and the upper term. The slowness of this splitting is the reason why for a long time the spectrum of helium was referred to as a doublet system.

We add a brief remark on the *notation for the terms*. A few lines above we have used capital letters  $S, P, D, \dots$  instead of small letters (p. 126). This is customary in the case of several electrons for the purpose of indicating the total orbital moment. It is also customary to define the multiplet character, that is to say, the number of terms belonging to a particular multiplet (all with the same principal and azimuthal quantum numbers), by attaching this number to the term symbol as a left-hand upper index; also, the multiplicity written down is always the one which occurs when  $l$  is large, viz.  $2s + 1$ . In fact, the vector sum  $\vec{l} + \vec{s} = \vec{j}$  gives the following values:

$$\begin{aligned} \text{for } l \geq s: & l - s, l - s + 1, \dots, l + s - 1, l + s \quad (2s + 1 \text{ values}), \\ \text{for } l \leq s: & s - l, s - l + 1, \dots, s + l - 1, s + l \quad (2l + 1 \text{ values}). \end{aligned}$$

If, for example,  $s = 1$ , we have a triplet; but this holds only from  $l = 1$  ( $P$ -term) upwards, while the term  $l = 0$  ( $S$ -term) is single. This peculiarity is ignored in the use of the notation; the two terms are written  $^3S, ^3P$ , although in reality the former of these is single.

The individual terms of the triplet  $^3P$  are distinguished by right-hand lower indices (suffixes), which indicate the value of  $j$ ; in helium, for instance,

$$\begin{aligned} & ^1P_1 \quad \text{for parhelium,} \\ & ^3P_0, ^3P_1, ^3P_2 \text{ for orthohelium.} \end{aligned}$$

Briefly, then, the notation used is as follows. The first number denotes the principal quantum number—it indicates the shell which the electrons occupy; if the electrons are in different shells, this quantum number must be omitted. The letter indicates the total orbital moment, the letters  $S, P, \dots$  corresponding to the quantum numbers  $l = 0, 1, \dots$ . The left-hand upper index gives the number of terms (at its maximum) in the term group defined by the principal quantum number and the letter—it is equal to  $2s + 1$ . The right-hand lower index distinguishes the individual terms of this group by assigning the inner quantum number  $j$  for the total angular momentum.

Returning once more to the term scheme for helium (p. 156), we have to point out that we have not shown in it the  $1^3S$ -term for orthohelium, i.e. the ground state to be expected. The reason is, that no lines have been found spectroscopically which could correspond to a transition to this term. We must conclude from this that in orthohelium the  $1^3S$ -state is missing. In fact, since we can calculate the position of this term approximately, we can tell approximately the

places in the spectrum at which the lines must lie which correspond to transitions to this term. In spite of the most thorough spectroscopic investigations in this region of the spectrum, not a single line has been found which can be brought into connexion with such transitions.

This result is quite inexplicable by the preceding principles as they stand. If we look for its cause in some special property of the term, we find the peculiarity that for this term all the quantum numbers are the same for the two electrons. Thus, both electrons lie in the  $K$  shell, and have therefore the same principal quantum number  $n_1 = n_2 = 1$ , as is also in fact the case for the parhelium term  $1^1S$ . Since the azimuthal quantum number of an electron can never be greater than  $n - 1$ , it must vanish in this case for both electrons:  $l_1 = l_2 = 0$ . (For higher  $S$ -terms, the two orbital moments could be different from zero, and by antiparallel setting cause the total orbital moment to vanish.) In orthohelium, the two electron spins are parallel, and the components of the spin moments are therefore also equal,  $\sigma_1 = \sigma_2 = \frac{1}{2}$ , so that the two sets of quantum numbers completely agree. In parhelium, however, the electron spins are antiparallel, so that the quantum numbers of the two electrons are separated in the projection of the two spin components.

### 5. Pauli's Exclusion Principle.

The state of matters just described in the case of helium, viz. the absence of the  $1^3S$ -term, the expected ground state of the orthohelium termsequence, suggested to Pauli (1925) a general examination of spectra, to see whether, in other elements and under other conditions, definite terms sometimes drop out. It was found that this actually happens; moreover, the term analysis showed in all cases that in these missing terms all the quantum numbers of the electrons agreed. Conversely, it was also found, the terms always drop out when these quantum numbers are the same. This discovery led Pauli to the following principle:

The quantum numbers of two (or more) electrons can never entirely agree; two systems of quantum numbers, which are deducible from each other by interchange of two electrons, represent *one* state.

The second part of the principle enunciates the *indistinguishability* of electrons. It is important for the enumeration of possible states, as required for the theory of the periodic system, and especially for statistics. It is to be remarked also that, in applying the principle, directional degeneration must be considered to be removed (for instance,

by an external magnetic field). The *individual* electron can then be characterized by *four quantum numbers*:

$$\left. \begin{aligned} n &= 1, 2, \dots, \\ l &= 0, 1, \dots, n-1, \\ j &= l - \frac{1}{2}, l + \frac{1}{2}, \\ m &= -j, -j+1, \dots, +j \end{aligned} \right\} \text{ or } \left\{ \begin{aligned} \mu &= -l, -l+1, \dots, +l, \\ \sigma &= -\frac{1}{2}, +\frac{1}{2}. \end{aligned} \right.$$

Here  $n$  denotes the principal quantum number, which can take all values from 1 upwards;  $l$  is the azimuthal quantum number, which runs from 0 to  $n-1$ ;  $j$  is the inner quantum number, and is only capable of the two values stated;  $m$  is the projection of  $j$  on the specially distinguished direction and, according to the rules for quantisation of direction, runs through the  $2j+1$  values between  $-j$  and  $+j$ . Alternatively, instead of  $j$  and  $m$  we may use the projections of  $l$

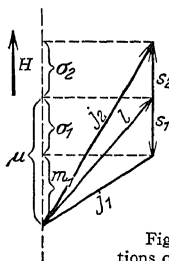


Fig 13.—Orbital moment  $l$ , spin moment  $s$ , total moment  $j$ ; with their projections on a specially distinguished direction.

and  $s$  on the special direction, viz.  $\mu$ , denoted previously by  $m_l$  (p. 145) and  $\sigma$  (see fig. 13).

Another concept which is frequently employed is this. Two electrons are said to be *equivalent* when they possess the same  $n$  and the same  $l$ . Two equivalent electrons must therefore, in accordance with Pauli's principle, differ from each other in the direction of  $l$ , or in the spin direction; and only certain definite values, not all values, of  $\mu$  and  $\sigma$  are possible for them. It is otherwise with two electrons which are not equivalent, but differ either in the principal or the azimuthal quantum number (or in both); in this case all values of  $\mu$  and  $\sigma$  are possible.

In Appendix XXIV (p. 319) we give an example of the enumeration of the terms for the two cases of non-equivalent and equivalent electrons. We shall not wait here to discuss these problems—which, though very simple in principle, are somewhat complicated to work out—but proceed at once to the most important and concrete application of Pauli's principle, the theory of the periodic system.

## 6. The Periodic System. Closed Shells.

We have already frequently referred to the periodic system, or table (see Table I, p. 35; fig. 14 shows a somewhat different form). The periodic table is an arrangement of the elements in a scheme



which originally was drawn up on the basis of their chemical behaviour and their atomic weights. It has turned out, however, that the real determining factor in the arrangement is not the atomic weight—consider isotopy—but the atomic number  $Z$ , i.e. the number of electrons revolving round the nucleus in the neutral atom. Now one of the most important applications of Pauli's principle is to the elucidation of the shell character of atoms, i.e. that property which finds its expression in the periodicity of the atoms with respect to chemical

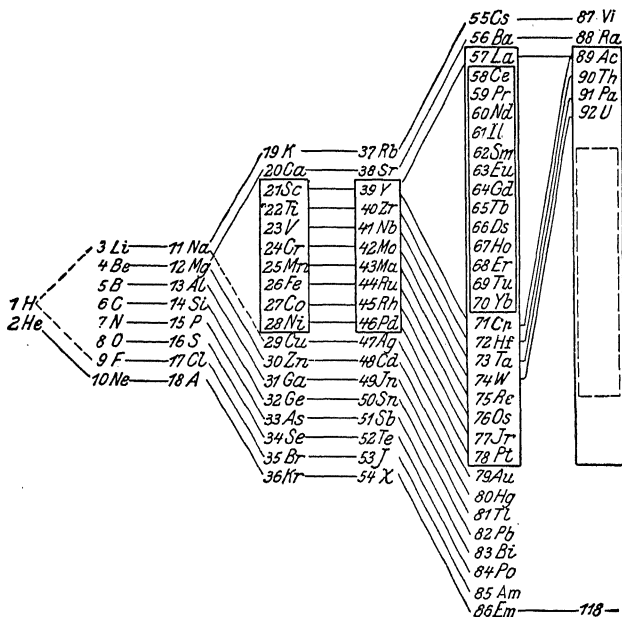


Fig. 14.—Diagram of the periodic system of the elements according to N. Bohr, exhibiting the synthesis of the shells, and the chemical relationship, of the atoms.

behaviour. We shall develop the theory of the periodic system by the method of proceeding step by step from a simpler element to the next higher. At each step, therefore, starting from an element with an electronic configuration which we know, we suppose the nuclear charge to be increased by a unit, and at the same time an electron to be inserted in the outer part of the electronic envelope of the known atom; what we want to know about is the way in which this addition of an electron takes place.

We begin with the simplest element—hydrogen ( $Z = 1$ ). Its electron in the ground state is in the lowest orbit, and has therefore for its

principal quantum number  $n = 1$ ; we say, as has already been remarked (p. 113), that the electron is in the  $K$  shell. With helium a second electron appears; thus for both these electrons  $n = 1$ , and therefore  $l = 0$ , and also  $\mu = 0$ . Hence, by Pauli's exclusion principle, they must possess different components of the spin moment. Since there are only two possible settings, viz.  $\sigma_1 = +\frac{1}{2}$  and  $\sigma_2 = -\frac{1}{2}$ , there are only places for two electrons in the  $K$  shell.

If we now add a third electron, there is no place for it in the  $K$  shell, and so it must settle in the  $L$  shell. Let us count the number of places in the  $L$  shell. First, as in the  $K$  shell, there are 2 electrons with  $l = 0$ , distinguishable only by their spin components; next come 6 places for electrons with  $l = 1$ , since of course  $\mu$  can take the three values  $-1, 0, 1$ , and for each of these there are two possible settings for the spin. There are therefore 8 places altogether in the  $L$  shell, and these are arranged (Stoner, 1924) in two sub-shells with 2 and 6 electrons. When the atomic number is steadily raised, the  $L$  shell becomes gradually filled; the atoms concerned are Li, Be, B, C, N, O, F, Ne. With Ne the  $L$  shell is completed; a new electron must settle in the  $M$  shell (Table IV, pp. 163, 164, 165).

We give the enumeration of the places in the  $M$  shell in the form of a table:

$l = 0, \mu =$	0,	$\sigma = \pm \frac{1}{2}:$	2 electrons
$l = 1, \mu =$	$-1, 0, 1,$	$\sigma = \pm \frac{1}{2}:$	6 ,,
$l = 2, \mu =$	$-2, -1, 0, 1, 2,$	$\sigma = \pm \frac{1}{2}:$	10 ,,
			in all 18 electrons.

In the  $M$  shell there are therefore 18 places, distributed over three sub-groups. The first two groups form together the second short period (of 8 elements); it extends from Na to A. Then, however, a deviation occurs from the previously invariable order of succession in which the electrons settle. This order is always determined by the energy released when a new electron settles (i.e. is bound). The energy relations, however, are not always such that one shell must be completed before an electron settles in the next shell. On the contrary, it may happen that an electron in an  $s$ -orbit of a higher shell is, from the energy point of view, more firmly bound than in a  $d$ - or  $f$ -orbit of the lower, still incomplete, shell. This case occurs in the further development of the  $M$  shell; the ten  $3d$ -terms for the 10 electrons still wanting are, as experiment shows, higher as regards energy than the  $4s$ -terms which correspond to binding of the electrons in the

TABLE IV  
DISTRIBUTION OF ELECTRONS IN THE ATOMS

Element		K		L		M			N		Ground Term (p. 158)	Ionization Potential (in electron volts)
		1, 0 1s		2, 0 2s	2, 1 2p	3, 0 3s	3, 1 3p	3, 2 3d	4, 0 4s	4, 1 4p		
H	1	1		—	—	—	—	—	—	—	$2S_{1/2}$	13.539
He	2	2		—	—	—	—	—	—	—	$1S_0$	24.45
Li	3	2		1	—	—	—	—	—	—	$2S_{1/2}$	5.37
Be	4	2		2	—	—	—	—	—	—	$1S_0$	9.48
B	5	2		2	1	—	—	—	—	—	$2P_{1/2}$	8.4
C	6	2		2	2	—	—	—	—	—	$3P_0$	11.217
N	7	2		2	3	—	—	—	—	—	$4S_{3/2}$	14.47
O	8	2		2	4	—	—	—	—	—	$3P_2$	13.56
F	9	2		2	5	—	—	—	—	—	$2P_{3/2}$	18.6
Ne	10	2		2	6	—	—	—	—	—	$1S_0$	21.5
Na	11	Neon Configuration				1	—	—	—	—	$2S_{1/2}$	5.12
Mg	12					2	—	—	—	—	$1S_0$	7.61
Al	13					2	1	—	—	—	$2P_{1/2}$	5.96
Si	14					2	2	—	—	—	$3P_0$	7.39
P	15					2	3	—	—	—	$4S_{3/2}$	10.3
S	16					2	4	—	—	—	$3P_2$	10.31
Cl	17					2	5	—	—	—	$2P_{3/2}$	12.96
A	18					2	6	—	—	—	$1S_0$	15.69
K	19	Argon Configuration						—	1	—	$2S_{1/2}$	4.32
Ca	20							—	2	—	$1S_0$	6.09
Sc	21							1	2	—	$2D_{3/2}$	6.57
Ti	22							2	2	—	$3F_2$	6.80
V	23							3	2	—	$4F_{3/2}$	6.76
Cr	24							5	1	—	$7S_3$	6.74
Mn	25							5	2	—	$6S_{5/2}$	7.40
Fe	26							6	2	—	$5D_4$	7.83
Co	27							7	2	—	$4F_{9/2}$	7.81
Ni	28							8	2	—	$3F_4$	7.606
Cu	29							10	1	—	$2S_{1/2}$	7.69
Zn	30							10	2	—	$1S_0$	9.35
Ga	31							10	2	1	$2P_{1/2}$	5.97
Ge	32							10	2	2	$3P_0$	7.85
As	33							10	2	3	$4S_{3/2}$	9.4
Se	34							10	2	4	$3P_2$	
Br	35							10	2	5	$2P_{3/2}$	11.80
Kr	36							10	2	6	$1S_0$	13.940

TABLE IV (*continued*)

Element	Configuration of Inner Shells	N		O			P	Ground-Term (p. 158)	Ionization Potential (in electron volts)
		4, 2 4d	4, 3 4f	5, 0 5s	5, 1 5p	5, 2 5d	6, 0 6s		
Rb 37	Krypton Configuration	—	—	1	—	—	—	$^2S_{1/2}$	4.16
Sr 38		—	—	2	—	—	—	$^1S_0$	5.67
Y 39		1	—	2	—	—	—	$^2D_{3/2}$	6.5
Zr 40		2	—	2	—	—	—	$^3F_2$	
Cb 41		4	—	1	—	—	—	$^6D_{1/2}$	
Mo 42		5	—	1	—	—	—	$^7S_3$	7.35
Ma* 43		6	—	1	—	—	—	$^6D_{9/2}$	
Ru 44		7	—	1	—	—	—	$^5F_5$	7.7
Rh 45	8	—	1	—	—	—	$^4F_{9/2}$	7.7	
Pd 46	10	—	—	—	—	—	$^1S_0$	8.5	
Ag 47	Palladium Configuration	—	—	1	—	—	—	$^2S_{1/2}$	7.54
Cd 48		—	—	2	—	—	—	$^1S_0$	8.95
In 49		—	—	2	1	—	—	$^2P_{1/2}$	5.76
Sn 50		—	—	2	2	—	—	$^3P_0$	7.37
Sb 51		—	—	2	3	—	—	$^4S_{3/2}$	8.5
Te 52		—	—	2	4	—	—	$^3P_2$	
I 53		—	—	2	5	—	—	$^2P_{3/2}$	10
Xe 54		—	—	2	6	—	—	$^1S_0$	12.078
Cs 55	The shells 1s to 4d contain 46 electrons	—	—	The shells 5s to 5p contain 8 electrons	—	1	$^2S_{1/2}$	3.88	
Ba 56		—	—		2	$^1S_0$	5.19		
La* 57		—	1		2	$^2D_{3/2}$			
Ce* 58		1	1		2	$^3H_4$			
Pr* 59		2	1		2	$^4K_{11/2}$			
Nd* 60		3	1		2	$^5L_6$			
Pr* 61		4	1		2	$^6L_{9/2}$			
Sm* 62		5	1		2	$^7K_4$			
Eu* 63		6	1		2	$^8H_{3/2}$			
Gd* 64		7	1		2	$^9D_2$			
Tb* 65		8	1		2	$^8H_{7/2}$			
Dy* 66		9	1		2	$^7K_{10}$			
Ho* 67		10	1		2	$^6L_{19/2}$			
Er* 68		11	1		2	$^5L_{10}$			
Tm* 69		12	1		2	$^4K_{17/2}$			
Yb* 70		13	1		2	$^3H_6$			
Lu* 71		14	1		2	$^2D_{3/2}$			

TABLE IV (*concluded*)

Element	Configuration of Inner Shells	O		P			Q	Ground Term (p. 158)	Ionization Potential (in electron volts)
		5, 2 5d	5, 3 5f	6, 0 6s	6, 1 6p	6, 2 6d	7, 0 7s		
Hf* 72	The shells 1s to 5p contain 68 electrons	2	—	2	—	—	—	$^3F_2$	
Ta* 73		3	—	2	—	—	—	$^4F_{3/2}$	
W 74		4	—	2	—	—	—	$^5D_0$	
Re* 75		5	—	2	—	—	—	$^6S_{5/2}$	
Os* 76		6	—	2	—	—	—	$^5D_4$	
Ir* 77	electrons	7	—	2	—	—	—	$^4F_{9/2}$	
Pt 78		8	—	2	—	—	—	$^3F_4$	
Au 79	The shells 1s to 5d contain 78 electrons	—	—	1	—	—	—	$^2S_{1/2}$	9.20
Hg 80		—	—	2	—	—	—	$^1S_0$	10.39
Tl 81		—	—	2	1	—	—	$^2P_{1/2}$	6.08
Pb 82		—	—	2	2	—	—	$^3P_0$	7.39
Bi 83		—	—	2	3	—	—	$^4S_{3/2}$	8.0
Po 84		—	—	2	4	—	—	$^3P_2$	
( )* 85		—	—	2	5	—	—	$^2P_{3/2}$	
Rn 86		—	—	2	6	—	—	$^1S_0$	10.689
( )* 87	Radon Configuration					—	1	$^2S_{1/2}$	
Ra* 88						—	2	$^1S_0$	
Ac* 89						1	2	$^2D_{3/2}$	
Th* 90						2	2	$^3F_2$	
Pa* 91						3	2	$^4F_{3/2}$	
U* 92						4	2	$^5D_0$	

The table shows not only those distributions of electrons and ground terms which have actually been ascertained from the spectra, but also those which have been determined by considerations of analogy; the latter are indicated by an asterisk attached to the symbol of the element. The probability of correctness of the inferences from analogy is, however, very high.

*N* shell. The next two electrons to be added (K, Ca) settle therefore in the *N* shell. Then, but not before, begins the filling up of the *M* shell, which continues from Sc ( $Z = 21$ ) to Zn ( $Z = 30$ ). The formation of the *N* shell then proceeds up to the element Kr ( $Z = 36$ ). The process described is repeated in the *N* shell; the two next electrons (Rb, Sr) settle in the *O* shell ( $n = 5$ ); the filling up of the *N* shell then begins, along with intermittent settlement of electrons in the *O* shell. The rare earths correspond to the completion of the *N* shell, which is accompanied by the binding of 8 or 9 electrons in the *O* shell and 2 in the *P* shell; this explains the great chemical similarity of these elements. All through, as has already been mentioned, the chemical behaviour of the elements is determined by the electronic configuration

in the outermost shell; elements with outermost shells of similar structure possess in large measure equivalent chemical properties. This is the explanation of the occurrence of periods in the system of the elements. Thus, for example, the inert gases Ne, Ar, Kr, Xe and Rn have all closed shells of 8 electrons; the alkalis are characterized by the fact that one electron revolves round the atom outside closed shells; the halogens, on the other hand, lack one electron to make up a closed shell. The problems of chemical binding and the chemical behaviour of the atoms will be taken up again in detail in Chapter VIII (p. 248).

## 7. Magnetism.

In § 2 (p. 140) we have investigated at length the splitting up of the terms of an atom in a magnetic field; it is determined by the directing force of the magnetic field acting on the magnetic moments of the spin and of the orbital motion of the electrons in the atom.

The magnetic moments within the atoms are also responsible for the magnetic behaviour of the substance made up of these atoms. Two cases have to be distinguished. A substance is called *paramagnetic* if its atoms (or molecules) possess a magnetic moment. The magnetization of the substance, which depends on this, is clearly in the same direction as the field, since of course it arises from the orientation of the magnetic moments in the direction of the field. Further, it is characteristically strongly dependent on the temperature, since the orientation of the elementary magnets is opposed by the smoothing out action upon direction which is due to the thermal motion. We shall deal with this more fully in Appendix XXV (p. 322). On the other hand, a substance is called *diamagnetic* if its atoms do not possess a permanent magnetic moment; in this case the magnetization which is produced is in the opposite direction to the field, and is practically independent of the temperature.

Atomic theory gives the magnetic moments (compare the results obtained above in connexion with the anomalous Zeeman effect, p. 140). Thus, for atoms with closed outer shells the moment 0 is found (in the inert gases, but also, according to measurements by Stern, in Zn, Cd, Hg). *S*-terms have no orbital moment ( $l = 0$ ), so that the magnetic moment of the atom is entirely due to the spin; thus the alkalis, with one radiating electron in an *s*-orbit, have a magnetic (spin) moment of 1 Bohr magneton, and the same is true also of the noble metals (Cu, Ag, Au). Those elements which have an intermediate shell not yet completed possess large magnetic moments (for instance, the elements of the iron group, and the rare earths).

The assertions of the theory about the magnitude of the magnetic moments of atoms can be tested by determining the magnetic susceptibility of the substances in question (Weiss, about 1910). However, this only gives a value averaged over all those directions, and all those values of magnetic moments, which occur; a direct test of the theory cannot therefore be immediately obtained by this method.

The magnitudes of the moments can also, however, be measured in the individual atom, as was shown by Stern and Gerlach (1921). Their method bears approximately the same relation to the macroscopic method of Weiss, as Aston's method for the determination of atomic masses bears to the macroscopic method of determining atomic weights, which only gives mean values over all the isotopes present. The method of Stern and Gerlach is based upon the deflection of a molecular beam in a non-homogeneous magnetic field. We may regard the atom with the magnetic moment as an elementary magnet of dimensions which though small are still finite. If we bring this magnet into a homogeneous magnetic field, it will move in a straight line in the field; for the magnetic force acts on its north pole with the same strength as on the south pole, but in exactly the opposite direction; thus the axis of the magnet may possibly execute a pendulum motion or precession about the direction of the field, but the centre of inertia of the magnet either remains at rest or moves in a straight line.

It is otherwise when the field is not homogeneous. In that case the forces acting on the north and south poles are not quite the same, so that, besides the couple which the two magnetic forces exert on the elementary magnet even in the case of a homogeneous field, a resultant force also acts on the magnet as a whole, and imparts to it an acceleration either in the direction of the magnetic field or in the opposite direction. If then an elementary magnet flies through a non-homogeneous magnetic field, it will be deflected from its rectilinear path. The amount of the deflection is determined by the degree of inhomogeneity of the field; in fact, to produce a deflection of sensible amount the inhomogeneity of the field must be so marked that the field changes decidedly even within the small length of the elementary magnet (which in our case is at most of the order of magnitude of atomic linear dimensions, viz.  $10^{-8}$  cm.). Stern succeeded in producing a sufficient degree of inhomogeneity by suitable construction of the pole-pieces of a magnet, one piece being shaped as a knife-edge, while the other piece, which he set opposite it, had a flat face, or was provided with a groove (fig. 15). The magnetic lines of force consequently crowd together at the knife-edge, so that the magnetic field strength

is considerably greater there than at the other pole-piece. A fine beam of atoms is now projected from an oven through a diaphragm system so as to pass between the pole-pieces. Each individual atom is deflected in the non-homogeneous field, according to the magnitude

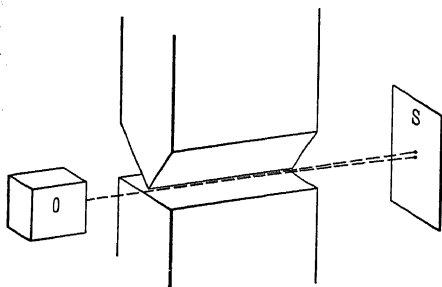


Fig. 15.—Diagrammatic representation of Stern and Gerlach's apparatus. A molecular beam issues from the oven *O*, and passes between the pole-pieces of the magnet (one of which has the form of a knife-edge) to the receiving screen *S*.

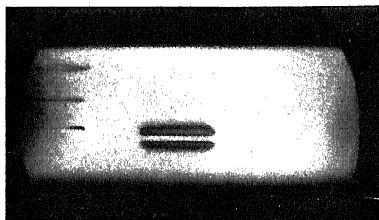
and direction of its moment. The traces of the individual atoms can be made visible on the receiving screen (by intensification if necessary, as in photography).

According to the classical theory, a broadening of the beam must be produced in this way on the screen, since the moments of the atoms flying through the magnetic field can, on that theory, have all directions relative to the field. On the quantum theory, however, on account of the quantising of direction, not all settings are possible, but only a certain discrete number, as we have seen above in detail when considering the anomalous Zeeman effect (p. 140). The beam, as it appears on the screen, will therefore be *split up* into a finite number of discrete beams; in fact, there must appear on the screen exactly  $2j + 1$  separate traces, if the atoms are in a state with the inner quantum number  $j$ ; for in this case there are just  $2j + 1$  possible settings, relative to the direction of the field, for the total angular momentum, and accordingly also for the total magnetic moment.

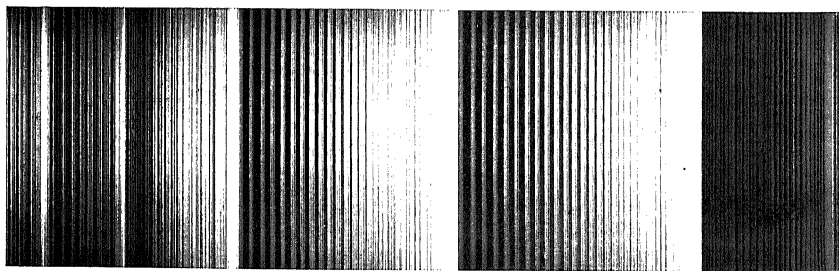
The experiment when carried out did in fact show a separation of the beam into several distinct beams; thus a beam of lithium atoms was split up into two beams (fig. 16, Plate VIII), as it should be, since the ground term of the lithium atom is a  $^2S$ -term (one valency electron in an  $s$ -orbit [ $l = 0$ ] with the spin moment  $\frac{1}{2}$ ). The magnitude of the magnetic moment also could be determined from the amount of the separation. In this way Gerlach succeeded in producing a direct proof that the magnetic spin moment is exactly equal to one Bohr magneton. A systematic investigation, upon various elements, yielded results which were throughout in complete agreement with the theory.

Later (1933), Stern succeeded in so refining the accuracy of the measurements as to make it possible to measure the magnetic moments





Ch. VI, Fig. 16.—Magnetic splitting up of a lithium beam by the method of Stern and Gerlach



Ch. VIII, Fig. 8.—Nitrogen bands (in the visible). Fine "triplets" are seen, which correspond to different rotational changes (by  $-1, 0, 1$ ); the regular sequence of these forms rotational bands with a "head". All these separate bands are associated with a jump of the vibrational quantum number, and with an electronic jump, which displaces the system into the visible.



of nuclei. In this apparatus, which is a marvel of refined workmanship, the atoms are detected by means of specially constructed, highly sensitive manometers, in which small changes of pressure are converted into thermal or electrical disturbances, and so made measurable. Interest, of course, is mainly directed to the proton. One is inclined to conjecture that its magnetic moment is given by the formula for the Bohr magneton, with the mass of the electron replaced by that of the proton. This *nuclear magneton* is therefore 1840 times smaller than the Bohr magneton. Experiments were made by Stern on molecular hydrogen, and by Rabi on atomic hydrogen; the difficulty in both cases is to reduce the magnetic moment of the outer electrons. The results agreed in giving the moment of the proton as two to three times greater than one nuclear magneton. We cannot therefore hope to obtain a serviceable theory of the proton from that of the electron (see § 8) simply by inserting the heavier mass instead of the lighter. The moment has also been measured for the heavy isotope *D* of hydrogen; it is less than a nuclear magneton.

### 8. Wave Theory of the Spin Electron.

In the preceding pages we have dealt with the theory of spin, and the questions connected with it—fine structure, Pauli's exclusion principle, &c.—solely on the basis of the vector model, regarding the angular momentum vectors as given magnitudes, and operating with them according to the rules of the classical or the Bohr theory, as the case might be. We mentioned at the outset (p. 146) that this procedure can be justified by wave mechanics. We cannot give a complete treatment here of the wave mechanics of the spinning electron, but we should like to point out at least in what way electronic spin is actually brought within the ambit of wave mechanics.

This extension of wave mechanics was introduced by Pauli (1925). The leading idea of Pauli's theory is somewhat as follows. For simplicity consider an isolated electron. Its state, according to Schrödinger, is described by a wave function  $\psi(x, y, z, t)$ , where  $|\psi|^2$  gives the probability of the electron being found at the point considered. We might now, keeping in view the idea of the rotating electron, introduce the spin into the wave equation by taking into account the rotational degrees of freedom. This, however, at once proves to be impossible; for in that case two new quantum numbers would appear in the solutions, as in every rotating system (e.g.  $l$  and  $m$  in the revolution of an electron); so that we would obtain an essentially greater number of states than the number actually found experimentally. This gives a new

proof that the idea of the rotating electron is not to be taken literally.

As has already been proved in § 1 (p. 136), it follows unequivocally from the spectra, especially those of the alkalis, that for a fixed set of three quantum numbers  $n, l, m$  the electron can be in two, and only two, different states, which have also different energies. We can take this new degree of freedom into account formally, by introducing besides the ordinary co-ordinates an additional co-ordinate  $\sigma$ , which can take only two values altogether; we shall denote these values by  $\sigma = +$  and  $\sigma = -$ , respectively. We can picture this to ourselves by supposing, say, that one value of this variable characterizes the state in which the spin is parallel to a specially distinguished direction, while the other value of the variable denotes the antiparallel setting. We thus obtain a wave function which now depends on five co-ordinates:  $\psi = \psi(x, y, z, t, \sigma)$ . It suggests itself, however, to split up this wave function into the two components

$$\psi = \begin{pmatrix} \psi_+(x, y, z, t) \\ \psi_-(x, y, z, t) \end{pmatrix}$$

which represent it respectively for the two possible values of the variable  $\sigma$ . It is evident that  $|\psi_+|^2$  specifies the probability of hitting upon the electron at the place considered, and that with a spin direction parallel to the special direction; while  $|\psi_-|^2$  is the corresponding probability for the opposite direction of the spin.

The question now meets us: how are we to calculate with these functions? For this Pauli has found the following method. In wave mechanics, there corresponds to every physical magnitude an operator, which is to be applied to the wave function. As such operators we can employ, as in Schrödinger's theory, differential operators (e.g. the component momenta  $p_x = (\hbar/2\pi i)(\partial/\partial x)$ ); or, as in Heisenberg's theory, matrices or other similar mathematical forms. It is therefore reasonable to associate operators also with the components  $s_x, s_y, s_z$  of the spin moment (in units  $\frac{1}{2}\hbar/2\pi$ ); these operators, however, do not act on the co-ordinates  $x, y, z, t$ , but on the two-valued variable  $\sigma$ . By this we are to understand that the application of these operators either changes the value of the magnitude  $\sigma$  or leaves it unchanged, while at the same time the whole wave function may possibly be multiplied by a factor. It has proved advantageous to regard the operators as "linear transformations" (matrices), and we usually employ for them the expressions

$$s_x = \begin{pmatrix} 0 & 1 \\ 1 & 0 \end{pmatrix}, \quad s_y = \begin{pmatrix} 0 & i \\ -i & 0 \end{pmatrix}, \quad s_z = \begin{pmatrix} 1 & 0 \\ 0 & -1 \end{pmatrix}.$$

If, generally,

$$a = \begin{pmatrix} a_{++} & a_{+-} \\ a_{-+} & a_{--} \end{pmatrix}$$

is such a matrix, then the operation  $a$ , applied to the vector of two components  $\psi = \begin{pmatrix} \psi_+ \\ \psi_- \end{pmatrix}$ , means the production of a new vector in accordance with the rule

$$a\psi = \begin{pmatrix} a_{++} & a_{+-} \\ a_{-+} & a_{--} \end{pmatrix} \begin{pmatrix} \psi_+ \\ \psi_- \end{pmatrix} = \begin{pmatrix} a_{++}\psi_+ + a_{+-}\psi_- \\ a_{-+}\psi_+ + a_{--}\psi_- \end{pmatrix}.$$

Thus we have

$$\begin{aligned} s_x\psi &= \begin{pmatrix} 0 & 1 \\ 1 & 0 \end{pmatrix} \begin{pmatrix} \psi_+ \\ \psi_- \end{pmatrix} = \begin{pmatrix} \psi_- \\ \psi_+ \end{pmatrix}, \\ s_y\psi &= \begin{pmatrix} 0 & i \\ -i & 0 \end{pmatrix} \begin{pmatrix} \psi_+ \\ \psi_- \end{pmatrix} = \begin{pmatrix} i\psi_- \\ -i\psi_+ \end{pmatrix}, \\ s_z\psi &= \begin{pmatrix} 1 & 0 \\ 0 & -1 \end{pmatrix} \begin{pmatrix} \psi_+ \\ \psi_- \end{pmatrix} = \begin{pmatrix} \psi_+ \\ -\psi_- \end{pmatrix}. \end{aligned}$$

The matrices are so chosen that they are subject to the same commutation rules as the ordinary components of angular momentum (§ 5, p. 127); further, in this case the  $z$ -direction is the specially distinguished direction, in which the spin, represented by the fourth co-ordinate, has set itself. For  $s_z$  is a diagonal matrix with the proper values (diagonal elements)  $+1, -1$ ; the  $z$ -component of the angular momentum has therefore one of the fixed values  $+1$  or  $-1$ . On the other hand,  $s_x$  and  $s_y$  are not diagonal matrices; their values are therefore not measurable simultaneously with  $s_z$ , but are only statistically determinate.

If now a magnetic field exists parallel to the  $z$ -axis, then, by § 2 (p. 144), the magnetic energy of the setting of the spin with respect to the field direction is given by

$$E_{\text{mag.}} = -2 \frac{e}{2\mu c} H \frac{1}{2} \frac{h}{2\pi} s \cos(sH) = - \frac{eh}{4\pi\mu c} H s_z.$$

The factor 2 here arises from the anomaly in the relation between magnetic and mechanical moment in the spin, as compared with the similar relation in the orbit—an anomaly mentioned on p. 142, and explained by Thomas as a consequence of the theory of relativity. In this equation  $s$  is the operator of the total spin moment, and  $s_z = s \cos(s, H)$  is its component in the direction of the field.

As for the wave equation itself, in the absence of magnetic fields it runs exactly like the Schrödinger equation for the problem in question:

$$(W_0 - E)\psi = 0.$$

Here  $W_0$  denotes the energy operator, consisting of the kinetic and the potential energy; for a one-electron problem we have

$$W_0 = \frac{1}{2m} \left( \frac{\hbar}{2\pi i} \right)^2 \Delta + V.$$

Since  $W_0$  does not act on the spin variable  $\sigma$ , we can split up the wave equation into the two components

$$(W_0 - E)\psi_+ = 0, \quad (W_0 - E)\psi_- = 0.$$

This signifies that without magnetic fields no transition is possible between the two possible settings of the spin, and that the electrons behave as if they possessed no magnetic moment.

If the atom is in a homogeneous magnetic field, the direction of which we shall take parallel to the  $z$ -axis, the usual energy operator  $W_0$  is supplemented by a term due to the magnetic field, viz.  $-(eh/4\pi\mu c)Hs_z$ , so that the wave equation now runs:

$$\left( W_0 - \frac{eh}{4\pi\mu c} Hs_z - E \right) \psi = 0.$$

Here again the separation into the two partial functions can be effected at once, taking into account the meaning of the operator  $s_z$ . We thus obtain the two equations

$$\left( W_0 - \frac{eh}{4\pi\mu c} H - E \right) \psi_+ = 0,$$

$$\left( W_0 + \frac{eh}{4\pi\mu c} H - E \right) \psi_- = 0,$$

from which we deduce immediately that the magnetic energy is simply added to the ordinary energy, so that an unperturbed term  $E_0$  is split up by the magnetic field into the two terms

$$E_0 + \frac{eh}{4\pi\mu c} H, \quad \text{and} \quad E_0 - \frac{eh}{4\pi\mu c} H.$$

This double possibility of setting is the reason for the occurrence of doublets in the spectra of the one-electron atoms.

We shall not proceed further here with the development of Pauli's

theory, but only remark that it has proved thoroughly successful, so long at least as the electron velocities are not too great. It must be emphasized, however, that Pauli's theory cannot be regarded as an explanation of the existence of spin, since of course when it was constructed the experimental facts, like the twofold possibility for the spin, and the ratio of the mechanical and magnetic moment, were simply inserted in the theory from the outside.

An essential advance in this direction is made in Dirac's *relativistic wave theory of spin* (1928). The object aimed at was to set up a wave equation which should satisfy Einstein's principle of relativity. According to this principle, the space co-ordinates  $x, y, z$  and the time  $t$  (the last multiplied by  $c\sqrt{-1}$ , where  $c$  is the velocity of light) stand on a precisely equal footing. But Schrödinger's differential equation p. 122) is not symmetrical in the four co-ordinates; it is of the second order in the spatial differential coefficients, but of the first order in the time derivative. Now Dirac has set up a wave equation which satisfies the relativistic postulate of symmetry, and, as regards form, is of the first order in all four variables. Although we cannot discuss this theory in detail, we have to emphasize particularly that, from very general principles, and with no special assumption about the spin, it deduces all those properties of the electron which we have summarized under the word spin (mechanical and magnetic moment in the correct ratio). In its results it is to a large extent equivalent to Pauli's theory, but it goes decidedly beyond the latter, especially in its conclusions regarding swift electrons. In respect of form, it differs from Pauli's theory in that it operates with four components of the wave function, these replacing Pauli's two functions  $\psi_+$  and  $\psi_-$ .

Although the conclusions of Dirac's theory agree very well, generally speaking, with experiment, yet even in it serious difficulties occur, which up till now have only in part been successfully overcome. There is in the first place the theory of the many-electron problem, which has not so far been successfully brought under the scheme of Dirac's theory; this is ultimately connected with the fact that even classically no satisfactory formulation of the two-body problem has yet been found within the ambit of the theory of relativity.

A further serious difficulty arises from the fact that in the theory of relativity the relation between the energy and the component momenta is quadratic in all these quantities, viz. (Appendix V, p. 271)

$$E^2 = G^2c^2 + m_0^2c^4.$$

For a given momentum  $G$  there are therefore always two values of the

energy,  $+E$  and  $-E$ . In the classical theory we accept only one of these values,  $+E$ , as having a physical meaning. This we have a right to do since doing so never leads to a contradiction. With the smallest value of  $G$ , viz.  $G = 0$ , we associate the value  $E_0 = +m_0c^2$ ; since  $G$  and  $E$  vary continuously, we have therefore, corresponding to all values of  $G > 0$ , the values  $E > m_0c^2 > 0$ . It is different in the quantum theory; here changes of state may take place by a jump; what is then to prevent  $E$  from falling by more than  $2m_0c^2$  in some process, and accordingly, by leaping over the excluded region from  $-m_0c^2$  to  $+m_0c^2$ , getting into the negative region  $E < -m_0c^2$ ? A simple example of this sort has been pointed out by O. Klein; it is called Klein's paradox. If an electron is caused to fly through two wire gratings, set up one behind the other (fig. 17), and a counter field is applied between the gratings, the electron in passing through loses part of its kinetic energy. If the retarding potential is increased, the loss of energy by the electron becomes greater and greater, until finally it is unable to pass through the gratings, but is reflected. So far, everything is classically in

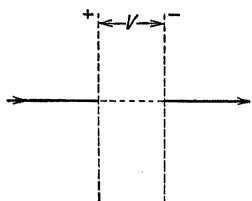


Fig. 17.—Illustration of Klein's paradox. According to Dirac's relativistic wave mechanics, an electron can overcome a counter field, if the potential difference exceeds twice the rest energy  $m_0c^2$  of the electron.

order. But if we go on to increase the counter field still further Dirac's theory shows that electrons can once more pass through the gratings, provided the potential difference between them is greater than twice the rest energy  $m_0c^2$  of the electron. In this case, however, the electrons emerge from the field with negative energy. If now we calculate the probability that the passage through this counter field will take place, we find that it increases with increasing field strength between the two gratings, but that for all electrostatic fields which can be produced experimentally it is vanishingly small. Although this is reassuring, and leaves unquestioned the applicability of Dirac's theory to ordinary problems, the theoretical difficulty in principle still remains.

Dirac has sought to solve the difficulty by the following bold line of argument.

In the formal theory, states of negative energy are on a perfectly equal footing with those of positive energy. This equality does not appear, however, to exist in nature; ordinary processes of motion all take place in the domain of positive energy. To give expression to



this fact, Dirac assumes that all the states of negative energy are normally occupied by negative electrons, which, however, do nothing to make themselves physically observable. We are therefore to consider that we are continually surrounded on all sides by electrons with all possible (negative) energies, without our noticing it at all. That of course implies very violent treatment of the ordinary accepted laws of electrodynamics; only electrons with positive energy can generate a field. From this sea of electrons, which extends from  $(-\infty)$  to the level  $-m_0c^2$ , electromagnetic fields, say in an act of absorption in the atom, can raise an electron, and bring it into the energy region  $E > m_0c^2$ ; it then takes its place as an ordinary electron and generates electromagnetic fields in accordance with the ordinary laws. There remains behind, however, an empty place in the sea of negative energy states—a hole—and clearly this has entirely the same character as a positive charge. Dirac's formulæ show that this hole in the midst of the ordinary negative electrons with negative energies behaves exactly like a positive electron with positive energy. Klein's paradox can now be explained very simply by the remark that the state of negative energy, into which the incident electron would go after passing through the opposing field, is already occupied by another electron; the process is therefore impossible.

Dirac thought at first that these holes corresponded to the protons, though he was conscious of the difficulty of the difference in the masses. But when the first indications of the existence of positrons were found by Anderson (see p. 43), Dirac at once saw in this a confirmation of his theory, and by prediction of phenomena took a hand in the direction of experimental research.

The most important consequence of the theory is that, by absorption of light quanta rich in energy, a pair, electron + positron, can be produced, and that conversely two such particles can unite, and produce a light quantum. The calculation of the probability for this follows the usual methods of wave mechanics, since the first process, according to Dirac, really consists in the raising of an electron out of a state (of negative energy), the second in the fall of an electron from a high state into a hole. Both processes are positively demonstrated to-day, as experimental results. We have the first in the showers which are produced by cosmic rays (see p. 43); further, cases are observed where by irradiation of light elements with  $\gamma$ -rays positive and negative electrons are liberated, which can be attributed to this process of materialization. In the converse process, the neutralization of two electrons of opposite sign, the energy  $2m_0c^2 = 10^6$  electron

volts becomes free (if the kinetic energy of the electrons is neglected); but momentum must also be conserved, and this excludes the case of emission of a light quantum, which of course has the definite momentum  $m_0c^2/c = m_0c$ . There are therefore two possible cases; either two light quanta originate in empty space, each equal to  $m_0c^2 = \frac{1}{2}10^6$  electron volts; or a particle in the neighbourhood takes up the residual momentum. If this is an electron, it will carry off a third of the energy  $2m_0c^2$ , so that the light quantum will retain  $\frac{2}{3}m_0c^2 = 680,000$  electron volts; if it is a nucleus, it will on account of its great mass take over very little energy, and the light quantum will therefore take the full amount of  $2m_0c^2 = 10^6$  electron volts. In point of fact, in scattering of  $\gamma$ -rays by atoms, besides the ordinary Compton scattering monochromatic waves are emitted, one of which has an energy of very exactly half a million electron volts; it arises presumably from the reunion of a positron, produced primarily at the absorption of the  $\gamma$ -rays, with a negative electron (Gray and Tarrant, Joliot). The process has been demonstrated unequivocally by Klemperer (1934). He made boron and carbon artificially radioactive by bombardment with protons and their heavy isotope (see § 5, p. 59); these materials then become a powerful source of positrons. The positrons are absorbed by a metal piece, which is enveloped on its two sides by two hemispherical counters, adjusted for  $\gamma$ -rays (see p. 32). It is then observed that the two counters give a simultaneous response, with a frequency exceeding that of any possible chance coincidences, so that two light quanta are being emitted at the same time in opposite directions; and it is found by absorption measurements that it is the half-million electron volts radiation which is involved. The other two processes are (in agreement with the theory) relatively extremely rare.

In spite of its successes, Dirac's hole theory of the positron is provisional rather than final. If a serious attempt is made to connect the theory with electrodynamics, while postulating that only free electrons and the holes, but not the great body of electrons of negative energy, shall act as generators of a field, the resulting formalism is extremely complicated, and seems far from satisfactory. Here theoretical physics is confronted with a serious problem.

## 9. Application of Wave Mechanics to the Nucleus.

The application of wave mechanics to the nucleus is met by the essential difficulty that the dimensions of the nucleus are of the same order of magnitude ( $10^{-13}$  cm.) as the diameter of the electron, which

TABLE V—NUCLEAR ANGULAR MOMENTUM  $I$ 

$Z$	Element	Isotopes	$I$
1	H	$\begin{cases} 1 \\ 2 \end{cases}$	$\begin{cases} \frac{1}{2} \\ 1 \end{cases}$
2	He	4	0
3	Li	$\begin{cases} 6 \\ 7 \end{cases}$	$\begin{cases} 0 \\ \frac{3}{2} \end{cases}$
7	N	14	1
8	O	16	0
9	F	19	$\frac{1}{2}$
15	P	31	$\frac{1}{2}$
17	Cl	35	$\frac{3}{2}$
25	Mn	55	$\frac{5}{2}$
29	Cu	63, 65	$\begin{cases} \frac{3}{2} \\ \frac{5}{2} \end{cases}$
48	Cd	$\begin{cases} 111, 113 \\ 110, 112, 114, 116 \end{cases}$	$\begin{cases} \frac{1}{2} \\ 0 \end{cases}$
80	Hg	$\begin{cases} 199 \\ 201 \\ 198, 200, 202, 204 \end{cases}$	$\begin{cases} \frac{1}{2} \\ \frac{3}{2} \\ 0 \end{cases}$
82	Pb	$\begin{cases} 207 \\ 204, 206, 208 \end{cases}$	$\begin{cases} \frac{1}{2} \\ 0 \end{cases}$
83	Bi	209	$\frac{9}{2}$

presumably represents a limit to the validity of the current theories (see § 2, p. 46). Nevertheless, we can make some confident affirmations about the quantum-mechanical properties of the nucleus, and by so doing bring various phenomena into rational relationship with each other. We might naturally expect that in general nuclei have angular momenta, which obey the laws of quantum mechanics (§ 5, p. 127). This is actually the case. The most important method for determining the quantum number  $I$ , which measures the angular momentum of the nucleus, depends upon the measurement of the hyperfine structure of spectral lines. With regard to some lines, as those of mercury, it had long been known that they consist of an aggregate of very fine lines; but it is only lately that these have been systematically investigated and brought into connexion with the nuclear moments. For the latter purpose, it must be assumed that the nucleus acts on the outer electrons not merely as a point charge, but with forces which depend on the orientation of its axis relative to the electrons. It is found that there are two kinds of hyperfine structure, one kind arising from the presence of several isotopic nuclei, each of which obviously

acts on the electrons a little differently from the rest; and the other kind occurring with a single isotope, and therefore dependent on direct, presumably magnetic coupling. The second kind is much the more important. The quantum-mechanical treatment is obvious. The nuclear angular momentum  $I$  and the electronic angular momentum  $J$  combine vectorially into the total moment  $F$ , and since each of the numbers is an integer or half an integer, definite configurations of lines arise, in the same way as in the Zeeman effect. By measurement of the separations and the intensities of the lines, we can determine  $I$ . We can also call in the aid of the Zeeman effect, and find out how the hyperfine structure changes in the magnetic field. In this way many  $I$ -values have been determined; a few of these are given in Table V. This shows that most of the elements with even atomic number show no splitting ( $I = 0$ ); the moments therefore seem to neutralize each other in pairs. Elements of odd atomic number have half-integral angular momenta, though nitrogen, it is true, has  $I = 1$ .

These results show clearly that in the nuclei the electrons are not to be counted as basic component parts (§ 4, p. 57). Consider the nitrogen nucleus, for example; if it were composed of 14 protons and 7 electrons, each of which has the moment  $\frac{1}{2}$ , the total moment would necessarily be half-integral; actually, however, it is 1. Conversely, the odd isotopes of Cd, Hg and Pb would have an odd number of electrons, and therefore an even number of particles altogether; but the total moment is half-integral. On the other hand, the facts are in agreement with Heisenberg's hypothesis of a structure consisting of protons and neutrons, each with moment  $\frac{1}{2}$ . The nitrogen nucleus would then consist of 7 neutrons and 7 protons, i.e. of 14 particles, in agreement with the evenness of the moment; and the same holds of Cd, Hg and Pb.

There is another method for the determination of the nuclear moment, based upon the intensity relations in the band spectra of molecules (§ 3, p. 238); but we cannot deal with this here.

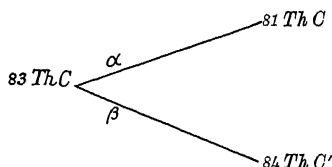
Closely connected with the mechanical moment of nuclei is their magnetic moment. We have already (§ 7, p. 166) made some remarks about the determination of this quantity, particularly for the proton. We may add here that it is also possible to evaluate magnetic nuclear moments by the theoretical calculation of the dimensions of the fine structure, and comparison with the measurements.

The quantities so far considered are properties of the nucleus in its natural state of equilibrium. But nuclei, as we have seen, can also be excited and disintegrated (§ 6, p. 61). The points for consideration are: what are the experimental facts, and how far can they be explained theoretically?

A general remark may be made in advance. The fact that we can make any attempt at all to study nuclear processes rests on our recognition of the circumstance that only heavy particles—neutrons, protons,  $\alpha$ -particles—come into account as elementary constituents of the nucleus, with no electrons or positrons. (We ignore  $\beta$ -disintegration here, relegating it to a special later investigation of the decomposition of neutrons into protons and electrons, or of protons into neutrons and positrons.) Light particles, in fact, would move so swiftly under the action of the nuclear forces that relativistic effects at their fullest intensity would require to be taken into account. A thoroughly satisfactory relativistic quantum mechanics is, however, not yet available—even allowing for Dirac's successes (§ 8, p. 169). With heavy particles, on the other hand, the velocities remain so small that we can apply the non-relativistic quantum mechanics as a sufficient approximation.

The first step is the experimental determination of energy levels in the nucleus. For this it would be necessary to excite nuclei to emit radiation, so that we might deduce the energy levels in the nucleus from the frequencies of this radiation, just as we do in the case of the external electrons, viz. by using Bohr's relation  $h\nu = E_2 - E_1$ . In point of fact, as we have already mentioned, in a few cases nuclei have successfully been caused by bombardment with  $\alpha$ -particles to emit  $\gamma$ -rays; but the material is still too meagre. More is got from the radioactive nuclei, which emit  $\gamma$ -rays spontaneously. The most fruitful method for determining their frequencies consists, not in this case in direct measurement, which on account of the shortness of the wavelengths is difficult, but in an indirect procedure, which, moreover, is of great interest for its own sake. We have already mentioned (§ 4, p. 30) that there are two kinds of  $\beta$ -emission: a continuous spectrum, attributed to-day to the breaking up of a neutron into proton + electron; and a discrete spectrum, of which we say, that it is of secondary origin. The evidence for the latter assumption is this: in the case of some substances, which emit  $\alpha$ - and  $\beta$ -rays simultaneously, we also find two transformation products (recognizable by their transformation constants). The bismuth isotope, ThC, splits up thus: 65 per cent with  $\beta$ -emission into the polonium isotope ThC', and 35 per cent with

$\alpha$ -emission into the thallium isotope  $\text{ThC}'$ , according to the scheme



In the case of other substances, however, e.g. radium and actinium, in spite of the most thorough investigation, the only product found is that corresponding to the  $\alpha$ -transformation. Consequently the  $\beta$ -rays cannot come out of the nucleus. Now some of these substances show a measurable  $\gamma$ -radiation; this suggests therefore that the  $\beta$ -rays come from the external electronic envelope, from which they are expelled by this  $\gamma$ -radiation.

The correctness of the suggestion is proved by the fact that simple relations exist between the energies of the  $\beta$ -rays and the energy levels of the electronic shells. In the first place we have

$$E_{\gamma} = E_{\beta} + A,$$

where  $E_{\gamma} = h\nu_{\gamma}$  is the energy of the  $\gamma$ -ray,  $E_{\beta}$  that of the  $\beta$ -ray, and  $A$  the energy required to remove the electron. Again, the same  $\gamma$ -ray can knock an electron out of the  $K$ ,  $L$  or  $M$  shell; between the energies of the  $\beta$ -rays we have then the relations

$$E_{\beta_1} + K = E_{\beta_2} + L = E_{\beta_3} + M = E_{\gamma}.$$

It has been found that the measured  $\beta$ -spectra do admit of grouping in accordance with these relations (Ellis (1921), Meitner). Moreover, it was possible to test the idea by causing the  $\gamma$ -rays from the lead isotope  $\text{ThB}$  to fall on lead, when it was found that the electrons produced photoelectrically had the same velocities as the  $\beta$ -rays from  $\text{ThB}$  itself. Further, observations on the electrons liberated from platinum by these  $\gamma$ -rays showed that their velocities were greater, by exactly the difference corresponding to the work of removal for lead and for platinum. Measurement of the  $\beta$ -energies then gave at the same time, by the preceding equation, the energy of the  $\gamma$ -ray. In this way it was found possible to draw up fairly extensive term schemes for various radioactive substances. We have therefore to picture radioactive disintegration in the following way (Meitner, 1925). The disintegration begins with  $\alpha$ -emission; the resulting nucleus is in an excited state, and reverts from this to the normal state with

$\gamma$ -emission. The proof of this rests on the high precision of the energy measurements, which allowed it to be decided that the rays come from the envelope of the final product, not from the original disintegrating substance.

The theory is capable of rendering a far-reaching account of this process of internal conversion of light energy into particle energy; the problem is that of finding how a spherical light wave, with its centre in the nucleus, is absorbed by the external electrons. The wave-mechanical calculation (Hulme, Taylor, Mott, 1932) furnishes results with regard to the relative intensities of the  $\beta$ -rays, originating in the  $K$ ,  $L$ ,  $M$  shells respectively. It is found, moreover, that besides the process of internal absorption of a  $\gamma$ -wave, a process is possible (Smekal, 1922) in which the excitation energy of the nucleus is, by means of the electromagnetic coupling, given up directly to the outside electrons, without a  $\gamma$ -ray appearing. By taking into consideration this radiationless process of conversion of nuclear energy into energy of the outside electrons (Fowler, 1930), we obtain better agreement of the calculated intensities with those observed. Another set of phenomena affording information about the energy levels of the nucleus is the fine structure of  $\alpha$ -rays (Rosenblum, 1931), and the occurrence of weak groups of  $\alpha$ -rays of abnormally long range (Rutherford, 1933). The explanation of these processes is not far to seek, if we remember that the  $\alpha$ -emission in these cases is accompanied by  $\gamma$ -emission. When the  $\alpha$ -ray is emitted, the residual nucleus can be left at various energy levels of short life; the energy of the  $\alpha$ -particles is therefore not unique, but has different values corresponding to the excitation levels of the nucleus. The nucleus subsequently returns to its ground state, with  $\gamma$ -emission. This is the explanation of the fine structure, which the use of powerful magnets has brought to light. But it can also happen, in an  $\alpha$ - or in a  $\beta$ -disintegration, that the resulting nucleus is left behind in various excitation levels of long life; in such a case the nucleus can either return to its ground state, with  $\gamma$ -emission, and then emit an  $\alpha$ -particle, or it can emit an  $\alpha$ -particle directly, with energy greater by the energy of the excitation level. In order that the latter event may happen often enough to be observable, the half-value period of the excited nucleus must be sufficiently short. But we know experimentally, and can also, as we shall see presently, explain theoretically that great  $\alpha$ -energy always runs parallel with shortness of life. We may thus understand the rare occurrence of  $\alpha$ -particles of very long range.

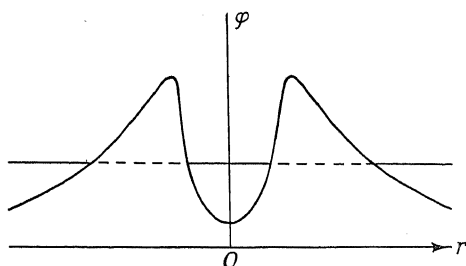
Here we come into contact with the fundamental law of radioactivity, the explanation of which is one of the greatest successes of quantum

mechanics. This law, which was discovered experimentally by Geiger and Nuttall (1911), gives a relation between the velocity  $v$  of the  $\alpha$ -particle and the life period  $T$  of the emitting nucleus: it runs

$$\log T = A + B \log v,$$

where  $A$  and  $B$  are constants.

The explanation, which was given independently by Gamow and by Condon and Gurney (1928), depends on a deep-seated distinction between quantum mechanics and ordinary mechanics, which is of importance in other cases also (as in cold electron emission, p. 222). In order to get a mechanical picture of the binding of an  $\alpha$ -particle to the rest of the nucleus, we must imagine a field of force which holds



it fast, and whose potential therefore has the form of a crater (fig. 18). We know the form of the outside walls of the crater, for the scattering experiments (§ 3, p. 51) tell us that Coulomb's

Fig. 18.—Crater form of the potential of the field of force which binds an  $\alpha$ -particle to the rest of the nucleus.

law holds down to very short distances—for uranium, certainly to less than  $r_0 = 3 \cdot 10^{-12}$  cm. The crater edge is therefore certainly higher than  $4e^2/r_0$  (where the factor 4 arises from the double charge of the  $\alpha$ -particle), and this is  $14 \cdot 10^{-8}$  erg. But the energy of the escaping  $\alpha$ -particles is less than half this value. If therefore an  $\alpha$ -particle is originally inside the crater, according to the laws of classical mechanics it could never get out.

It is different in quantum mechanics. Here the motion of the particle within the crater is represented by a wave—a standing wave, in fact. Consider an optical analogy, the total reflection of a ray of light, which passes through glass and meets at a sharp angle the plane surface separating glass and air. From the corpuscular point of view no light penetrates into the air. According to the wave theory, however, there exists in the air also a sort of wave disturbance, which, it is true, carries off no energy, and only penetrates a few wave-lengths into the air space. But if we now set a second piece of glass with its face parallel to the face of the other, and at such a short distance that the disturbance in the air gap reaches it with intensity not too much



impaired, then a small amount of energy at once passes into the second body, and the incident wave is propagated through it, though of course very much weakened. The weakening obviously diminishes rapidly, as the air gap is made narrower, until in the limit when the two pieces of glass come into contact, the ray goes through with intensity unchanged.

In wave mechanics, we have quite a similar state of affairs at the walls of our crater. If these are of finite height and thickness, the standing waves inside the crater are by no means completely obliterated at the wall, but are propagated in it in weakened form, and emerge at the outside as progressive waves. Since, however, the square of the amplitude,  $|\psi|^2$ , gives the probability of occurrence of the particles, we must conclude that the particles can penetrate through the wall with a frequency which is given by  $|\psi|^2$ . This will clearly diminish extremely rapidly with increasing thickness of the rampart; here the important matter is the thickness at the height which corresponds to the energy of the  $\alpha$ -particles (see fig. 18, p. 182). Swift  $\alpha$ -particles will therefore have a much greater probability of getting out than slow ones; this corresponds qualitatively to Geiger and Nuttall's law. The exact analysis (Appendix XXVI, p. 326) leads to a somewhat complicated formula of similar character; to obtain this, detailed assumptions about the interior of the energy crater are not required—its radius at the upper edge is really all that matters. For this radius a value in the neighbourhood of  $8 \cdot 10^{-13}$  cm. is found in all radioactive nuclei, and it is just this constancy of the effective radius which forms the foundation of Geiger and Nuttall's formula, with fixed constants  $A$  and  $B$ . The penetrability of the potential wall explains the fact (§ 6, p. 61) that the disintegration of lithium nuclei by protons can be observed for so small velocities that the distance of nearest approach is much larger than the radius of the Li nucleus. The proton wave penetrates into the interior of the Li nucleus and determines a small probability for the presence of a proton inside the nucleus, which becomes unstable and breaks into two  $\alpha$ -particles.

The model used in these considerations is highly simplified; only one particle is assumed to move, the other constituents of the nucleus not being considered separately, but only as the sources of the force acting on the moving particle. Recently various attempts have been made to treat the nucleus as a system of many particles with the help of the methods of non-relativistic quantum mechanics. For this purpose definite assumptions about the forces between the elementary particles (neutrons and protons) are necessary. It then becomes pos-

sible to determine the stationary states of the nucleus and compare the energy levels with the observations mentioned above (p. 179); the most important is, of course, the ground state the energy of which gives the mass of the nucleus. The same forces come into play when a beam of particles (protons, neutrons) is scattered by a nucleus; the observation of scattering processes gives therefore an independent test of the assumptions made.

We know the Coulomb repulsion between two protons. But this force is certainly small compared with those forces of unknown origin which keep the particles together; their range must be of the order of the nuclear radius ( $10^{-13}$  cm.). Wigner (1933) assumed these to be ordinary central forces falling off with the distance according to a law  $Ae^{-r/a}$ . Heisenberg (1932) proposed a theory where the forces were similar to the "exchange" forces by which quantum mechanics explains chemical valency (see p. 253) and which depend on the parallel or antiparallel setting of the spins of the particles involved. A modification of Heisenberg's assumptions about the spins has been introduced by Majorana. Both these theories are able to account for the empirical fact that the binding energy varies as  $N$  and the nuclear radius as  $\sqrt[3]{N}$ , if  $N$  is the number of protons and neutrons in the nucleus. A law containing two constants, like that mentioned above, explains roughly the behaviour of the masses. But we shall not enter into a discussion of the results obtained, as the fundamental conceptions are rather doubtful.

Bohr (1936) has shown that the experiments on absorption and scattering of neutrons by nuclei can be qualitatively explained if one remembers that the nucleus consists of many particles all interacting with forces of the same order, so that it is not possible to talk about excited states of a single particle. He even thinks it likely that the protons and neutrons in the nucleus entirely lose their individuality.

A complete theory of the nucleus will doubtless only be found in connexion with an explanation of the real nature of the elementary particles, and their relation to electromagnetic fields.

## CHAPTER VII

### Quantum Statistics

#### 1. Heat Radiation and Planck's Law.

We have already mentioned several times that the quantum theory is of statistical origin; it was devised by Planck when he was endeavouring to deduce the law of heat radiation (1900).

If, contrary to the order of historical development, we have discussed the quantum theory of the atom before quantum statistics, we have our reasons. In the first place, the failure of the classical theory displays itself in atomic mechanics—for instance, in the explanation of line spectra, or the diffraction of electrons—even more immediately than in the attempts to fit the law of radiation into the frame of classical physics. In the second place, it is an advantage to understand the mechanism of the individual particles and the elementary processes before proceeding to set up a system of statistics based upon the quantum idea.

In this chapter we shall supply what is lacking, and begin with the deduction of the *law of heat radiation*, following Planck's method. We think of an enclosure, say a box, whose walls are heated by some contrivance to a definite temperature  $T$ . The walls of the enclosure send out energy to each other in the form of heat radiation, so that within the enclosure there is a radiation field. We characterize this electromagnetic field by specifying the energy density  $u$ , which in the case of equilibrium is the same for every internal point; if we split up the radiation into its spectral components, we denote by  $u_\nu d\nu$  the energy density of all radiation components whose frequency falls in the interval between  $\nu$  and  $\nu + d\nu$ . Thus the function  $u_\nu$  extends over all frequencies from 0 to  $\infty$ ; it represents a continuous spectrum. Up till now we have been occupied with line spectra, which are emitted by individual atoms in rarefied gases. But even molecules, consisting of only a limited number of atoms, do not send out isolated lines as single atoms do, but narrow "bands", often not resolvable. And the more numerous and more densely packed the atoms are, the

more the lines run into one another so as to form continuous strips. A solid represents an infinite number of vibrating systems of all frequencies, and therefore emits a continuous spectrum.

Now there is a theorem of Kirchhoff's (1859) which states that the ratio of the emissive and absorptive powers of a body depends only on the temperature of the body, and not on its nature; otherwise radiative equilibrium could not exist within a cavity containing substances of different kinds. (By emissive power is meant the radiant energy emitted by the body per unit time, by absorptive power the fraction which the body absorbs of the radiant energy which falls upon it.) By a *black body* is meant a body with absorptive power equal to unity, i.e. a body which absorbs the whole of the radiant energy falling upon it. The radiation emitted by such a body—called “black radiation”—is therefore a function of the temperature alone, and it is important to know the spectral distribution of the intensity of this radiation. The following pages are devoted to the determination of the law of this intensity.

With regard to the experimental production of black radiation, it has been proved by Kirchhoff that an enclosure (oven) at uniform temperature, in the wall of which there is a small opening, behaves as a black body. In fact, all the radiation which falls on the opening from the outside passes through it into the enclosure, and is there, after repeated reflection at the walls, finally completely absorbed by them. The radiation which emerges again from the opening must therefore possess exactly the spectral distribution of intensity which is characteristic of the radiation of a black body. The same also holds good of the radiation in the interior, for the radiation emerging from the opening is of course simply the totality of the components of the cavity radiation, which fall upon the opening from within.

Without going beyond thermodynamics and the electromagnetic theory of light, we can deduce two laws regarding the way in which black body radiation (or, as it is also called, cavity radiation) depends on the temperature. *Stefan's law* (1879) states that the total emitted radiation is proportional to the fourth power of the temperature of the radiator; the hotter the body, the more it radiates. Proceeding a step further, W. Wien found the *displacement law* (1893) which bears his name, and which states that the spectral distribution of the energy density is given by an equation of the form

$$u_\nu = \nu^3 F\left(\frac{\nu}{T}\right),$$

where the function  $F$  is a function of the ratio of the frequency to the temperature, but cannot be defined more precisely from the standpoint of thermodynamics. These two theorems can be proved, as is shown in Appendix XXVII (p. 329) by treating the radiation as a thermodynamic engine which—by means of a hypothetical movable mirror—can do work in virtue of the radiation pressure; so that, on account of the Doppler effect, arising from the motion of the mirror, the frequency of the radiation is altered, and therefore its energy content also. We may remark that Wien's law includes Stefan's law; to deduce the latter, we have only to integrate over the whole spectrum:

$$\int u_\nu d\nu = \int \nu^3 F\left(\frac{\nu}{T}\right) d\nu.$$

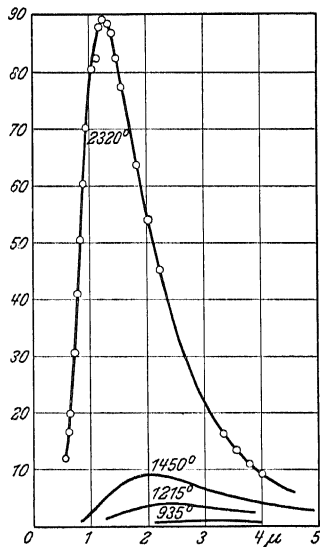
Putting  $x = \nu/T$ , and taking  $x$  as a new variable of integration, we find

$$\int u_\nu d\nu = T^4 \int x^3 F(x) dx,$$

so that the total radiation energy is proportional to the fourth power of  $T$ , since the integral with respect to  $x$  is independent of  $T$ , being a mere constant.

The reason for calling Wien's law the "displacement law" is this. It was found experimentally that the intensity of the radiation from an incandescent body, maintained at a definite temperature, was represented graphically, as a function of the wave-length, by a curve of the form sketched in fig. 1. For extremely short, as also for extremely long waves, the intensity is vanishingly small; hence it must have a maximum value for some definite

Fig. 1.—Distribution of intensity of heat radiation as a function of the wave-length, according to measurements by Lummer and Pringsheim.



wave-length  $\lambda_{\max}$ . If we now change the temperature of the radiating body, the intensity graph also changes; in particular the position of the maximum is shifted. In this way it was found from the measurements that the product of the temperature and the wave-

length, for the corresponding maximum of intensity, is constant; or

$$\lambda_{\max} T = \text{const.}$$

This relation is explained at once by Wien's law. In the first place, the energy density and the radiation intensity have the same spectral distribution; they differ, as can be shown from theoretical considerations, only by the constant factor  $c/4$ . Further, we have always up to now referred to the energy distribution as a function of the frequency  $\nu$ ,  $u_\nu$  representing the radiation energy in the frequency interval  $d\nu$ . The displacement law, however, refers to a graph showing the intensity distribution as a function of  $\lambda$ , so that we have now  $u_\lambda$  representing the energy in the wave-length interval  $d\lambda$ . The conversion from  $u_\nu$  to  $u_\lambda$  is easy: obviously we must have  $u_\nu d\nu = u_\lambda d\lambda$ ; and, since  $\lambda\nu = c$ , we have, as the relation between  $d\nu$  and  $d\lambda$ ,  $|d\nu|/\nu = |d\lambda|/\lambda$ . Hence, for the spectral distribution of energy expressed as a function of the wave-length, we find

$$u_\lambda = \frac{c^4}{\lambda^5} F\left(\frac{c}{\lambda T}\right).$$

We can now prove the displacement law at once, by calculating the wave-length for which  $u_\lambda$  is a maximum. The condition for this is  $du_\lambda/d\lambda = 0$ , or

$$\frac{c^4}{\lambda^5} \left\{ -\frac{5}{\lambda} F\left(\frac{c}{\lambda T}\right) - \frac{c}{\lambda^2 T} F'\left(\frac{c}{\lambda T}\right) \right\} = 0,$$

from which there follows

$$\frac{c}{\lambda T} F'\left(\frac{c}{\lambda T}\right) + 5F\left(\frac{c}{\lambda T}\right) = 0.$$

This is an equation in the single variable  $c/\lambda T$ , whose solution, assuming it exists, must of course have the form

$$\lambda T = \text{const.}$$

Thus the theorem about the displacement of the intensity maximum with temperature follows immediately from Wien's law. The value of the constant, it is true, cannot be determined until the special form of the function  $F$  is known.

About this function, however, thermodynamics by itself has nothing to say; to determine it, we must have recourse to special represen-

tations by a model. All the same, it is clear from thermodynamical considerations that the form of the law given by the function  $F$  must be independent of the special mechanism. As the simplest model of a radiating body, Planck therefore chose a linear harmonic oscillator of proper frequency  $\nu$ . For this oscillator, we can on the one hand determine the energy radiated per second; this being the radiation emitted by an oscillating dipole (p. 133), it is given by

$$\delta\epsilon = \frac{2e^2\overline{\dot{r}}^2}{3c^3} = \frac{2e^2}{3mc^3} (2\pi\nu)^2 \bar{\epsilon},$$

where  $\epsilon$  is the energy of the oscillator, and the bars denote mean values over times which, while great in comparison with the period, are yet sufficiently small to allow us to neglect the emitted during its time of continuance. For the oscillator we course

$$\frac{m}{2} \overline{\dot{r}^2} = \bar{\epsilon}_{\text{kin}} = \bar{\epsilon}_{\text{pot}} = \frac{1}{2} \bar{\epsilon}.$$

On the other hand, the work done on the oscillator per second by a radiation field with the spectral energy density  $u_\nu$  is

$$\delta W = \frac{\pi e^2}{3m} u_\nu,$$

as follows from the equation of motion of the oscillator, and is proved in Appendix XXVIII (p. 333). In the case of equilibrium, these two amounts of energy must of course be equal. Hence we have

$$u_\nu = \frac{8\pi\nu^2}{c^3} \bar{\epsilon}.$$

If therefore we know the mean energy of an oscillator, we know also the spectral intensity distribution of the cavity radiation.

The value of  $\bar{\epsilon}$ , as determined by the methods of classical statistics (§ 6, p. 9), would be

$$\bar{\epsilon} = kT,$$

where  $k$  is Boltzmann's constant; for, according to Boltzmann's theorem (proved at p. 14), in the state of equilibrium the value  $\epsilon$  for the energy of the oscillator occurs with the relative probability  $e^{-\epsilon/kT}$ , so that we obtain  $\bar{\epsilon}$  by averaging over all values of  $\epsilon$ , with this weight factor. Putting for brevity  $\beta = 1/kT$ , we find

$$\bar{\epsilon} = \frac{\int_0^{\infty} \epsilon e^{-\beta \epsilon} d\epsilon}{\int_0^{\infty} e^{-\beta \epsilon} d\epsilon} = -\frac{d}{d\beta} \log \int_0^{\infty} e^{-\beta \epsilon} d\epsilon$$

$$= -\frac{\partial}{\partial \beta} \log \frac{1}{\beta} = \frac{1}{\beta} = kT.$$

This mean value  $\bar{\epsilon}$  is independent of the nature of the radiating system (oscillator), as follows at once from the method of calculating it; it must hold in exactly the same form for a system of atoms or other radiating particles.

If the classical mean value of the energy of the oscillator, as thus determined, is substituted in the radiation formula, the result is

$$u_{\nu} = \frac{8\pi\nu^2}{c^3} kT.$$

This is the *Rayleigh-Jeans radiation law* (1900, 1909). We see in the first place that it agrees, as of course it must, with Wien's displacement law, which, as a deduction from thermodynamics, holds in all cases. For the long-wave components of the radiation, i.e. for small values of  $\nu$ , it also reproduces the experimental intensity distribution very well; in this region the intensity of the radiation increases with the frequency, in proportion to its square. For high frequencies, however, the formula fails; we know from experiment that the intensity function reaches a maximum at a definite frequency and then decreases again. The above formula, however, fails entirely to show this maximum; on the contrary, according to the formula, the spectral intensity distribution increases as the square of the frequency, and for extremely great frequencies, i.e. for extremely short waves, becomes infinite; the same is true of the total energy of radiation  $u = \int_0^{\infty} u_{\nu} d\nu$ ; the integral diverges. We have here what is called the "ultra-violet catastrophe".

Attempts have been made to remedy this very conspicuous failure in the theory by the following hypothesis. Suppose that, as in chemistry, a finite reaction time is needed before equilibrium is attained, and let the reaction velocity in the case of cavity radiation be very small, so that a very long time must elapse before equilibrium is reached, during which time the system as a whole, in consequence of external influences, will in general, we may well imagine, have completely changed. But this certainly does not go to the root of the matter;



for we can suppose of course, at least theoretically, that the cavity can be maintained unchanged at constant temperature for as long as we please, so that the above abnormal state of equilibrium would necessarily set in some time after all.

In these circumstances, Planck put forward the bold idea that these difficulties can be removed by postulating *discrete, finite quanta of energy*  $\epsilon_0$ ; the energy of the oscillators is to be (besides  $\epsilon = 0$ ) either equal to  $\epsilon_0$  or to  $2\epsilon_0$  or to  $3\epsilon_0$ , and so on. In this way we do in fact obtain Planck's radiation law, which is verified experimentally with striking accuracy. The essential point is the determination of the mean energy  $\bar{\epsilon}$ ; in point of form the investigation only differs from the earlier one (p. 190) in replacing integrals by sums; for the individual energy values, now as before, occur with the frequency given by the Boltzmann factor, but now only the energy amounts  $n\epsilon_0$  ( $n = 0, 1, 2, 3, \dots$ ) are possible, instead of all amounts as before. Hence, we have for the mean value,

$$\begin{aligned}\bar{\epsilon} &= \frac{\sum_{n=0}^{\infty} n\epsilon_0 e^{-\beta n\epsilon_0}}{\sum_{n=0}^{\infty} e^{-\beta n\epsilon_0}} = -\frac{d}{d\beta} \log \sum_{n=0}^{\infty} e^{-\beta n\epsilon_0} \\ &= -\frac{d}{d\beta} \log \frac{1}{1 - e^{-\beta\epsilon_0}} = \frac{\epsilon_0 e^{-\beta\epsilon_0}}{1 - e^{-\beta\epsilon_0}} \\ &= \frac{\epsilon_0}{e^{\beta\epsilon_0} - 1}, \quad (\beta = 1/kT)\end{aligned}$$

so that, on substitution of this expression in the radiation formula, we obtain

$$u_\nu = \frac{8\pi\nu^2}{c^3} \frac{\epsilon_0}{e^{\epsilon_0/kT} - 1}.$$

In order that this formula may not be inconsistent with Wien's displacement law, which, being deduced from thermodynamics alone, is certainly valid, we must assume that

$$\epsilon_0 = h\nu,$$

where  $h$  is a universal constant (Planck's constant); for the temperature can only appear in the formula in the combination  $\nu/T$ . This gives *Planck's radiation law*

$$u_\nu = \frac{8\pi h\nu^3}{c^3} \frac{1}{e^{h\nu/kT} - 1}.$$

As we remarked above, this radiation formula is in very good agreement with experimental results. The distribution of intensity for various temperatures is shown graphically in fig. 2. For low frequencies the function increases approximately as the square of  $\nu$ ; for, if  $h\nu/kT \ll 1$ , we can expand the exponential function in the denominator, and obtain a series in ascending power of  $h\nu/kT$ , the first

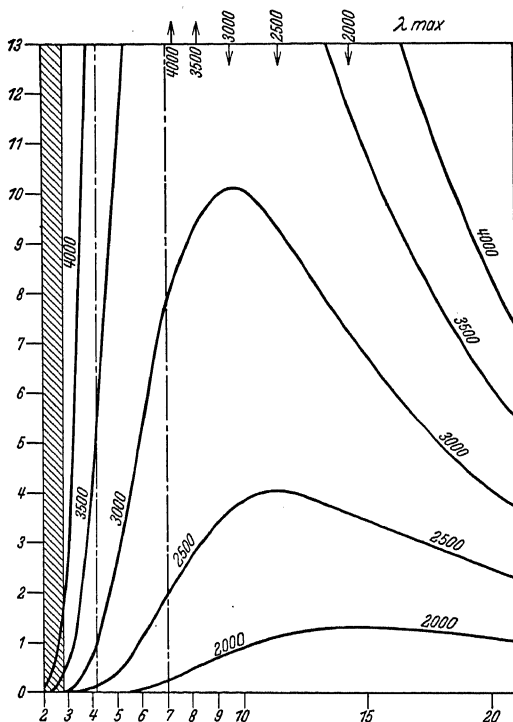


Fig. 2.—Spectral distribution of intensity of heat radiation according to Planck, for temperatures from 2000° to 4000°. The unit for the abscissæ is  $0.1\mu$ . The shaded part represents the ultra-violet region, as far as  $0.28\mu$ ; the visible region lies between the two dotted lines. The numbers on the curves are absolute temperatures. Note the agreement with the experimental results shown in fig. 1, p. 186.

term of which corresponds precisely to the Rayleigh-Jeans radiation formula:

$$u_\nu = \frac{8\pi h\nu^3}{c^3} \frac{1}{(1 + h\nu/kT + \dots) - 1} = \frac{8\pi\nu^3}{c^3} kT + \dots;$$

for long-wave radiation Planck's formula therefore agrees with the classical one. Not so in the short-wave region; if  $h\nu/kT \gg 1$ , the

exponential function has a value much greater than unity, so that we can leave out the 1 in the denominator; for short waves we thus obtain the radiation law in the approximate form

$$u_\nu = \frac{8\pi h \nu^3}{c^3} e^{-h\nu/kT},$$

a form which had previously (1896) been obtained by Wien experimentally, from his measurements in this region of the spectrum (fig. 3). Between the domains of these two limiting laws there lies a continuous transitional region, which contains the maximum of the distribution curve. This maximum, in accordance with Wien's law, is displaced towards the region of short waves as the temperature rises. Calculation of the intensity maximum leads in the same way as before to the relation

$$\frac{kT}{h\nu_{\max}} = \frac{k}{hc} \lambda_{\max} T = C,$$

where  $C$  is a constant, which is found by solving a transcendental equation, and has the value 0.2014.

We have mentioned above (Chap. I, § 8, p. 23) that

the atomic constants  $k$  and  $h$  can be determined by means of the spectral distribution of energy in black body radiation. The first determination of  $h$ , in fact, was effected by Planck by means of Stefan's law and Wien's displacement law. According to the former, the total energy radiated, from 1 sq. cm. of the surface of a body heated to temperature  $T$ , is  $J = \sigma T^4$ ;  $\sigma$  is called the radiation constant, and has the value  $\sigma = 5.77 \times 10^{-5}$  erg/(cm.<sup>2</sup> sec. degree<sup>4</sup>). We obtain it theoretically by integrating Planck's distribution function over the

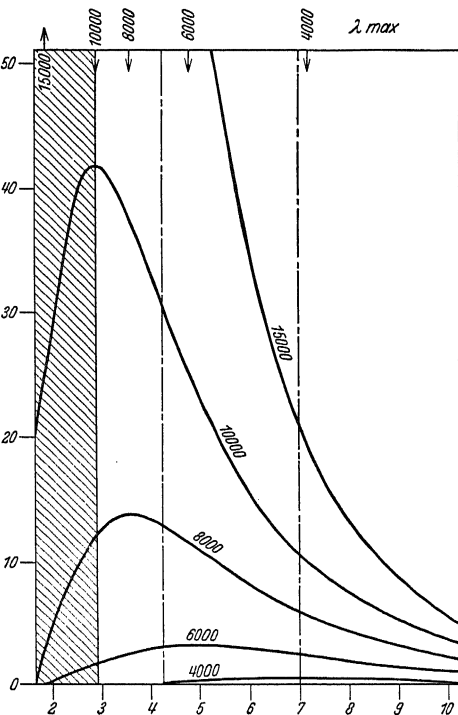


Fig. 3.—The same as fig. 2 for higher temperatures; the unit for the ordinates is smaller than in fig. 2

whole spectrum. On the other hand, the measurements give for the constant in Wien's law ( $\lambda_{\max}T = \text{const.}$ ) the value 0.288 cm. degree. (To give an idea of its order of magnitude, it is sufficient to remark that the intensity maximum of the radiation from the sun, which radiates like a black body at the temperature  $T = 6000^\circ$ , lies in the green region of the spectrum, i.e. at a wave-length of about 4500 Å.) From these two empirically determined constants,  $h$  and  $k$  can be calculated; the values so obtained agree very well with the correct ones.

We may add one brief remark. From the general course of the distribution curves we can see how low the efficiency is, when glowing heated bodies, such as incandescent lamps, are used for purposes of illumination; for the visible spectral region occupies only a small strip in the figure giving the intensity curves for the heat radiation, and the rest of the radiated energy is lost, so far as illumination is concerned.

To return to the method of proving Planck's formula: it goes without saying that his hypothesis met at first with the most violent opposition. Physicists were unwilling to believe that the law of radiation could only be deduced by introducing the quantum hypothesis; this they regarded as a mathematical artifice, which they considered might in some way or other be interpreted without going outside the circle of classical ideas. All attempts, however, at such an interpretation have been a complete failure.

To Einstein is due the credit of having been the first to point out that, apart from the law of heat radiation, there are other phenomena which can be explained by means of the quantum hypothesis, but which are unintelligible from the standpoint of classical physics. In the year 1905, Einstein put forward the *light quantum hypothesis*, and as experimental evidence for it cited the law of the photoelectric effect; this we have already dealt with (§ 2, p. 69).

## 2. Specific Heat of Solids and of Polyatomic Gases.

A year or two later (1907) Einstein showed that Planck's formula for the mean energy of an oscillator,

$$\bar{\epsilon} = \frac{h\nu}{e^{h\nu/kT} - 1},$$

finds direct confirmation in the thermal behaviour of solids. It is a fact of experiment that at high temperatures a law holds called *Dulong and Petit's law*, which states that the molecular heat, i.e. the specific

heat per mole (p. 8), is approximately 6 cal/degree for all solids. From the classical standpoint this law is at once intelligible. In a solid, every atom can be regarded as a harmonic three-dimensional oscillator, since of course in a solid we think of the atom as quasi-elastically bound to a certain position of equilibrium. The mean kinetic energy falling to them as three-dimensional oscillators, by the rules of classical statistics, is  $3kT$ ; so that a mole of the substance possesses the energy  $U = 3LkT = 3RT$ , where  $R$  is the gas constant, and is approximately equal to 2 cal/degree. The specific heat is found from this as the increase of energy for a temperature increase of  $1^\circ$ , so that

$$c_v = \frac{\partial U}{\partial T} = 3R \sim 6 \text{ cal/degree.}$$

Experiment, however, shows deviations from this rule; and the harder the body, i.e. the more firmly the atoms are bound to their equilibrium positions, the greater are the deviations; thus, for diamond, the specific heat per mole at room temperature is only about 1 cal/degree.

Einstein's explanation of these deviations is that it is not permissible in this case to use the classical expression for the mean energy of the oscillators, but that we must apply the expression obtained by Planck for the mean energy of a quantised oscillator. In that case the mean energy of the oscillators per mole is

$$U = \frac{3Lh\nu}{e^{h\nu/kT} - 1} = 3RT \frac{h\nu/kT}{e^{h\nu/kT} - 1}.$$

In this formula  $h\nu$  is the elementary quantum of the vibrational energy of the oscillators; it is so much the greater, the more firmly the atoms are bound to their equilibrium positions; loose binding is equivalent to small vibrational energy, and so to low frequency. The important point now is, whether  $h\nu$  is less or greater than  $kT$ . In general, at room temperature  $h\nu/kT \ll 1$ , so that we can simplify the formula for the energy of the oscillator by expansion. Hence in this case it passes over into the classical formula

$$U = 3RT \frac{h\nu/kT}{(1 + h\nu/kT + \dots) - 1} = 3RT + \dots,$$

and so gives Dulong and Petit's law.

If, however, the atoms are very firmly bound (as in diamond), or if the specific heat is measured at very low temperatures,  $h\nu/kT$  becomes comparable with unity, or even greater than unity, and deviations occur from Dulong and Petit's law. We then obtain for

the specific heat a curve of the form shown in fig. 4, which for large  $T$  rises asymptotically to the classical value of 6 cal/degree, falls away for low temperatures, and at  $T = 0$  passes through the origin. Experimental investigations to test the results given by this theory, mainly carried out by Nernst and his colleagues, showed approximate agreement between theory and experiment, particularly

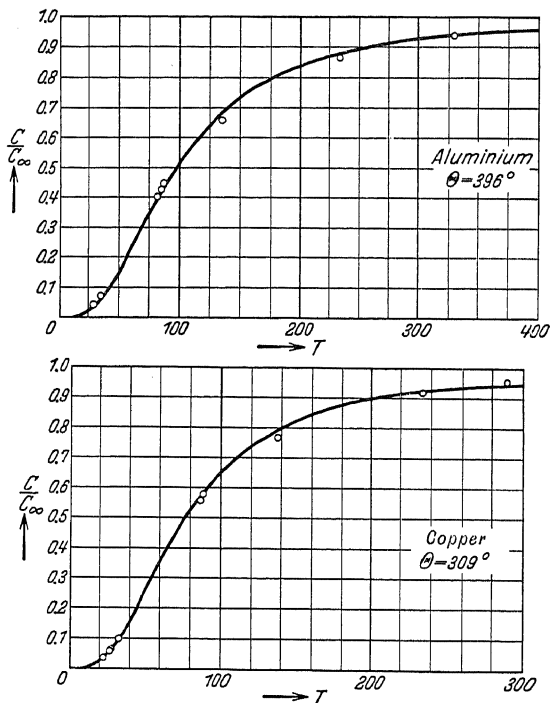


Fig. 4.—Graph of specific heats at low temperatures according to Debye; the small circles show observed points, the continuous curves correspond to Debye's theory.  $\theta$  is a temperature characteristic of the substance, such that  $C(=C_v)$  is a function of  $T/\theta$ .

as regards the fact that the specific heat approaches zero as the temperature diminishes. Nevertheless they indicated deviations which proved that the theory in the form in which it then existed still required improvement.

These improvements were made by Debye, and independently by Born and Kármán (1912). They rest on the following considerations. Up to this point we have dealt with the individual atoms of the solid (crystal), as if they performed undisturbed harmonic vibrations independently of one another. This, however, is by no means the case,

since of course the atoms in the lattice are very strongly coupled together. We therefore ought not to say that the  $L$  atoms in the crystal perform oscillations of the same frequency, but must rather deal with the coupled system of  $3L$  vibrations (corresponding to the  $3L$  degrees of freedom of the  $L$  atoms per mole), and accordingly write down the energy in the form

$$U = \sum_{r=1}^{3L} \frac{h\nu_r}{e^{h\nu_r/kT} - 1},$$

where  $\nu_r$  is the frequency corresponding to an individual vibration.

It would of course be a difficult undertaking to try to evaluate this sum directly on the basis of some assumed model. But we can obtain approximate formulæ; the simplest method—which serves moderately well for monatomic crystals—was given by Debye, and is as follows. The proper vibrations of the individual atoms of the crystal lattice appear in the ordinary theory of elasticity as elastic proper vibrations of the whole crystal; though, it is true, only waves whose lengths are much greater than atomic distances actually come under observation (as sound waves). Hence, for an approximate evaluation of the above energy sum, we may replace the spectrum of the proper vibrations of the individual atoms by the spectrum of the elastic vibrations of the whole crystal. We have therefore to solve the problem of determining the spectrum of the elastic vibrations of a body which from the standpoint of the theory of elasticity is regarded as continuous. Since similar problems (*counting proper vibrations*) will also occupy us later, we shall deal with this one here in some detail.

For simplicity, we take a body in the form of a cube, of side  $a$ , and try to find its proper vibrations, that is to say, the standing waves

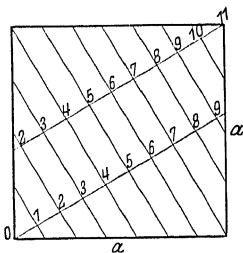


Fig. 5.—Example of a proper vibration in a cubical enclosure; every side must correspond to a whole number of half wave-lengths. (In the example  $k_1 = 7$ ,  $k_2 = 4$ .)

which can exist in this cube. This is the same problem, in three dimensions, as we have already met with in the one-dimensional case (vibrating string); and we also mentioned the two-dimensional case of the vibrating circular membrane (p. 125). For every proper vibration it is necessary that each side of the cube should cut off an exact whole number of half wave-lengths (fig. 5). If therefore a plane wave is propagated in this cube, in the direction specified by the

direction cosines  $\alpha, \beta, \gamma$  (so that  $\alpha^2 + \beta^2 + \gamma^2 = 1$ ), the projection of a side on this direction must be an integral multiple of  $\lambda/2$ . Thus we obtain the three equations

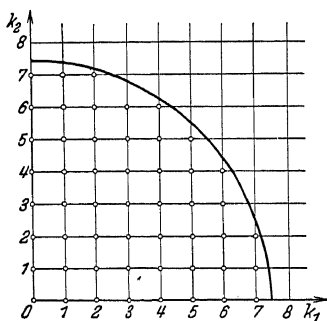
$$k_1 \frac{\lambda}{2} = \alpha a, \quad k_2 \frac{\lambda}{2} = \alpha \beta, \quad k_3 \frac{\lambda}{2} = \alpha \gamma.$$

The case  $\alpha = 1, \beta = \gamma = 0$  is trivial; the wave advances parallel to the first side, and the proper vibrations are the same as in the case of a string fixed at both ends. But the foregoing conditions are necessary for the existence of a proper vibration, in the three-dimensional case also.

We can now consider the spectrum of these proper vibrations. In the first place, the three conditions, by squaring and adding, give

$$k_1^2 + k_2^2 + k_3^2 = \frac{4a^2}{\lambda^2},$$

so that the wave-lengths of the proper vibrations are determined by means of the sums of the squares of three whole numbers  $k_1, k_2, k_3$ .



We may with advantage represent the individual proper vibrations by a point in the space of the  $k$ 's, by associating the point in  $k$ -space whose three co-ordinates are  $k_1, k_2, k_3$  with the proper vibration defined by the three preceding equations. But the number of proper vibrations with a wave-length  $> \lambda$  is

Fig. 6.—Counting proper vibrations: the number of vibrations with wave-length  $> \lambda$  is equal to the number of lattice points inside the circular arc shown, and this is approximately equal to the area of the quadrant.

then just the number of points  $(k_1, k_2, k_3)$  which lie within a sphere of radius  $2a/\lambda$  and in the first octant of the  $k$ -space (because the foregoing equation  $k_1^2 + k_2^2 + k_3^2 = 4a^2/\lambda^2$  represents a sphere in the  $k$ -space, with centre at the origin and radius  $2a/\lambda$ ). All proper vibrations with greater wave-length than  $\lambda$  are therefore represented by points which lie within this sphere in  $k$ -space. The limitation to the first octant of the sphere is obviously needed, since of course the integers  $k$  are positive.

The required number of points is approximately equal to the volume of the spherical octant just defined. We may easily convince ourselves of this with the aid of fig. 6 (though this, of course, is only two-dimensional). Since the co-ordinates of the lattice-points are whole numbers, the small squares formed by the lattice-lines have the



area 1. The error made in replacing the number of lattice-points within the circle (sphere) by its area (volume) is small compared to the correct value itself, if the number of points within the circle (sphere) is large. The number of proper vibrations  $> \lambda$  is therefore equal to the volume of the spherical octant, i.e.

$$Z = \frac{1}{8} \frac{4\pi}{3} \left( \frac{2a}{\lambda} \right)^3 = \frac{4\pi}{3} \frac{V}{c^3} \nu^3,$$

where  $c$  denotes in this case the velocity of the elastic sound waves. From this we get at once the *number of proper vibrations in the frequency interval between  $\nu$  and  $\nu + d\nu$* :

$$dZ = \frac{4\pi V}{c^3} \nu^2 d\nu.$$

We shall return to this formula later; the method of proving it is general, and not confined merely to elastic proper vibrations. Moreover, as has been shown by Weyl, the formula is true whatever may be the form of the volume  $V$ .

We now return to the determination of the mean energy of the atoms in a solid. The formula already deduced (p. 197), viz.

$$U = \sum_{r=1}^{3L} \frac{h\nu_r}{e^{h\nu_r/kT} - 1},$$

can now be transformed by the method used above, i.e. by regarding the frequencies  $\nu_r$  no longer as proper vibrations of the atoms, but as frequencies of elastic waves in the body. Since we know the number of these vibrations in the frequency interval between  $\nu$  and  $\nu + d\nu$ , we can rewrite the sum as an integral, in the integrand of which this number of vibrations appears as the weighting factor for this vibrational state, alongside the original summand; thus, roughly

$$U = \int \frac{h\nu}{e^{h\nu/kT} - 1} \frac{4\pi V}{c^3} \nu^2 d\nu.$$

It must be observed, however, that the total number of proper vibrations of the crystal is finite, viz. equal to  $3L$ ; a maximum frequency  $\nu_m$  therefore exists, which is defined by the equation

$$3L = Z = \frac{4\pi V}{3c^3} \nu_m^3,$$

or

$$\nu_m = c \sqrt[3]{\left( \frac{9L}{4\pi V} \right)}.$$

The preceding integral is therefore not to be taken to  $\infty$ , but only to this limiting frequency as its upper limit. For the energy  $U$  we thus obtain the expression

$$\begin{aligned}
 U &= \frac{4\pi V}{c^3} \int_0^{\nu_m} \frac{h\nu^3 d\nu}{e^{h\nu/kT} - 1} = \frac{4\pi V}{c^3} \left(\frac{kT}{h}\right)^4 h \int_0^{x_m} \frac{x^3 dx}{e^x - 1} \\
 &= 3RT \frac{3}{x_m^3} \int_0^{x_m} \frac{x^3 dx}{e^x - 1},
 \end{aligned}$$

where for brevity  $x_m$  is written for  $h\nu_m/kT$ . This relation is known as *Debye's formula*, which clearly gives a better approximation than Einstein's model of independent oscillators, all with the same frequency (fig. 4, p. 196). Debye's formula was still further improved by Born and Kármán, who in the enumeration of the proper vibrations took into account the lattice structure of crystals, to which in the preceding counting process we have paid no heed whatever. They thus obtained separate formulæ for the various types of crystal structure.

Empirically, this refined theory has been confirmed throughout. E.g. experiment gives a  $T^3$ -law for the specific heat at extremely low temperatures, whereas the simple Einstein theory leads to an exponential law for its increase with temperature. Debye's theory, however, leads to the correct law for low temperatures; in fact, for these temperatures,  $x_m$  tends to infinity, so that the integral in the formula for  $U$  becomes practically constant, while the fourth power of  $T$  stands before the integral; the specific heat is found from the energy by differentiation with respect to the temperature, and the experimental  $T^3$ -law therefore follows as it ought to do. Recent investigations by Blackman (1935), it is true, have shown that sometimes the measured specific heat seems to obey the  $T^3$ -law at temperatures not low enough to justify the preceding theoretical explanation. He has shown that a refined application of the theory of Born and v. Kármán which takes account of the lattice structure of the crystals is able to explain these cases; according to the arguments he adduces there should be an increase of specific heat at still lower temperatures, and the real  $T^3$ -law should only be expected at the very lowest temperatures, not yet employed for this measurement. This increase of the specific heat is actually observed, and many other predictions of the theory are also verified. One might say that at the present time the experimental facts relating to the specific heat of solids have been completely explained.

The quantum theory of the oscillator can be applied in an exactly

similar way to *polyatomic gases*. Here also the specific heat, as determined experimentally, increases with temperature in accordance with Planck's formula corresponding to the molecular vibrations. We have seen in Chap. I (§ 5, p. 9) that the classical theory of specific heats has to assume that the electronic motion in the atom does not contribute to specific heat, but that the heat content comes from the motions of the atoms and molecules themselves; the atoms having only translatory motions, while diatomic molecules have two rotational degrees of freedom (round the two axes perpendicular to the neutral line). The explanation of this fact, which was rather puzzling to pre-quantum physics, is of course the enormous size of the binding energy of the electrons compared with the average energy  $kT$  of thermal motion. If we replaced the electrons by virtual oscillators having frequencies  $\nu$  corresponding to the terms of the line spectra, then the energies  $h\nu$  would be large compared with  $kT$ . At normal, and even very high temperatures, none of these oscillators would be excited. This explains immediately why they do not contribute to the specific heat.

Similar considerations hold for the rotations of molecules consisting of light atoms. We have seen that the rotational energy is given by  $E_j = j^2 \hbar^2 / 8\pi^2 A$ ,  $j = 0, 1, 2, \dots$ , corresponding to Bohr's quantum theory (p. 100); wave mechanics replaces  $j^2$  by  $j(j+1)$  (p. 129), but in both cases the difference between neighbouring energy levels is  $\epsilon = \hbar^2 / 4\pi^2 A$ . Here  $A$  means the moment of inertia of the molecule, and  $\epsilon$  is greatest when  $A$  is smallest, i.e. for the molecule  $H_2$ . One should expect that for temperatures given by  $kT < \epsilon$  the rotational energy of  $H_2$  should vanish; this is really the case, as was first observed by Eucken (1912), who found that at about  $40^\circ$  K. the specific heat of  $H_2$  has decreased to the value for a monatomic gas. For all other gases the critical temperature is too low to allow an observation of this effect of the complete "freezing in" of the molecular rotations. But the beginning of the process, a decrease of the specific heat, can be observed, and from this the moment of inertia calculated. The results are in good agreement with other observations, as on band spectra (p. 238), and with calculations from theoretical models.

### 3. Quantisation of Cavity Radiation.

We now return to the law of cavity radiation. We have seen in the preceding section that Planck's hypothesis has been brilliantly successful not only for cavity radiation, but also in the theory of specific heats. The latter success furnishes additional strong support for the quantum theory.

On the other hand, with regard to the deduction of the radiation law by Planck's method, it is unsatisfactory, in so far as it is based in part on the laws of classical physics and only in part appeals to the quantum theory. The method by which the formula is obtained, connecting the mean energy of an oscillator with the radiation field in a cavity, viz.  $u_\nu = (8\pi\nu^2/c^3)\bar{\epsilon}$ , is purely classical; in deducing it, the classical laws of absorption and emission by an oscillator are employed.

Can this circuitous route via the absorbing and emitting oscillator be avoided? The idea suggests itself, in view of the methods of Rayleigh-Jeans and Debye, that we should try to deal with the electromagnetic field, within a cubical cavity with reflecting walls, in exactly the same statistical way as with the proper vibrations of crystals in the theory of specific heats. The cavity with reflecting walls possesses precisely the same kind of proper vibrations as a crystal, and we can therefore calculate the number of these vibrations in a definite frequency interval between  $\nu$  and  $\nu + d\nu$  in the same way as before:

$$dz = 2 \frac{4\pi V}{c^3} \nu^2 d\nu.$$

We have here, however, an extra factor 2, since for each possible wave-length and direction of wave normal there are two different waves, corresponding to the two independent directions of polarization.

It is now only consistent to assume that every proper vibration behaves like a Planck oscillator with the mean energy

$$\frac{h\nu}{e^{h\nu/kT} - 1},$$

where we are simply extending the method of calculation and line of thought of last section, from the proper vibrations of the crystal lattice to the proper vibrations of the cavity. In this way we again obtain Planck's formula

$$\begin{aligned} u_\nu d\nu &= \frac{1}{V} dz \frac{h\nu}{e^{h\nu/kT} - 1} \\ &= \frac{8\pi h\nu^3}{c^3} \frac{1}{e^{h\nu/kT} - 1} d\nu. \end{aligned}$$

Although this method of deduction is formally extremely simple, it contains all the same a serious difficulty of principle. The formula used for the mean energy of an oscillator is bound up with the idea

that an oscillator of frequency  $\nu$  can possess not only the energy  $h\nu$ , but also integral multiples of this energy quantum, the frequency of occurrence of the energy  $n h\nu$  in an oscillator being proportional to  $e^{-nh\nu/kT}$ . Thus we obtained the mean energy of an oscillator as the mean value

$$\frac{\sum_0^{\infty} n h\nu e^{-nh\nu/kT}}{\sum_0^{\infty} e^{-nh\nu/kT}} = \frac{h\nu}{e^{h\nu/kT} - 1}.$$

If we extend these ideas to the proper vibrations of the cavity, treating those as if they were oscillators, it follows that an electromagnetic proper vibration of frequency  $\nu$  can have the quantity of energy  $n h\nu$ . This would mean, in the language of the light quantum theory, that a light quantum (photon), which is associated with an electromagnetic vibration of frequency  $\nu$ , can take the energy values  $n h\nu$ ; this is contrary to Einstein's light quantum hypothesis, according to which the energy of a light quantum is always  $h\nu$ . This hypothesis of Einstein's, however, has to-day, on the evidence of experiments relating to the elementary atomic processes, reached the status of an assured fact. The attempt to explain the law of cavity radiation by means of classical statistics, merely supplemented by the quantum hypothesis, therefore fails when we try to translate the deduction of the radiation formula by undulatory considerations into the language of the corpuscular theory.

#### 4. Bose-Einstein Statistics of Light Quanta.

We can also try to deduce the radiation formula, not as above from the pure wave standpoint by quantisation of the cavity radiation, but from the standpoint of the theory of light quanta, that is to say, of a corpuscular theory. For this we must therefore develop the statistics of the light-quantum gas, and the obvious suggestion is to apply the methods of the classical Boltzmann statistics, as in the kinetic theory of gases; the quantum hypothesis, introduced by Planck in his treatment of cavity radiation by the wave method, is of course taken care of from the first in the present case, in virtue of the fact that we are dealing with light quanta, that is, with particles (photons) with energy  $h\nu$  and momentum  $h\nu/c$ . It turns out, however, that the attempt to deduce Planck's radiation law on these lines also fails, as we proceed to explain.

We do not in this case characterize a radiation component, as we

id before, by its wave-length and the direction cosines of its wave normal, but by the component momenta associated with it in de Broglie's theory, viz.

$$p_x = \frac{h}{\lambda} \alpha, \quad p_y = \frac{h}{\lambda} \beta, \quad p_z = \frac{h}{\lambda} \gamma.$$

Since, however, as we know from our previous investigations (p. 198), the only standing waves which can exist in a cubical cavity (of side  $a$ ) are those which satisfy the conditions

$$k_1 = \frac{2a}{\lambda} \alpha, \quad k_2 = \frac{2a}{\lambda} \beta, \quad k_3 = \frac{2a}{\lambda} \gamma,$$

it follows that the *momentum space* of the light quanta is also *discontinuous*; in fact the two sets of equations give

$$p_x = \frac{h}{2a} k_1, \quad p_y = \frac{h}{2a} k_2, \quad p_z = \frac{h}{2a} k_3;$$

we therefore see that the component momenta can only be integral multiples of  $h/2a$ . The momentum itself is given by

$$p = \frac{h}{\lambda} = \frac{h\nu}{c} = \frac{h}{2a} \sqrt{(k_1^2 + k_2^2 + k_3^2)}.$$

In the classical statistics (§ 6, p. 9) the momentum space was divided into cells of arbitrary form and magnitudes  $\omega_1, \omega_2, \dots$ ; the ratio of the size of the cell to the total momentum space  $\omega$ , i.e.  $g_s = \omega_s/\omega$ , was called the statistical weight of the cell. Here the momentum space is covered discontinuously; the statistical weight of a cell  $\omega_s$  is therefore simply the number of points  $(k_1, k_2, k_3)$  which lie within it. This follows from the principle of correspondence, as it is necessary to have a continuous transition to classical theory. There is therefore a smallest cell, among those containing any point at all; it is customary in quantum statistics to use the word "cell" for this smallest cell, i.e. really for the triplet of numbers  $k_1, k_2, k_3$ . We must therefore find another name for the elementary region of statistics  $\omega_s$ . We shall call  $\omega_s$  a "sheet". The number of cells  $(k_1, k_2, k_3)$  in the sheet  $\omega_s$  is then its statistical weight  $g_s$ . We now choose for the form of the sheet the spherical shell between the momenta  $p$  and  $p + \Delta p$ . Being merely a formal expedient, the division into sheets is quite arbitrary. The number of cells (e.g. the number of possible states) that is contained in such a sheet is

found from the formula deduced above (p. 202) for the number of proper vibrations of the cavity in the frequency interval between  $\nu$  and  $\nu + \Delta\nu$ , if we remember the relation  $p = h\nu/c$ :

$$g_s = 2 \frac{4\pi V}{c^3} \nu^2 \Delta\nu = 2 \frac{4\pi V}{h^3} p^2 \Delta p.$$

We can now see, with this expression before us, that the result could also be obtained by dividing up the whole phase space—i.e. co-ordinate space and momentum space together—into smallest cells of magnitude  $h^3$ , and then counting the number of these cells which lie in that region of phase space which corresponds to the spatial volume  $V$  and the region between  $p$  and  $p + \Delta p$  in the momentum space; taking account of the doubling of states due to polarization, we thus find

$$g_s = 2 \frac{4\pi p^2 \Delta p V}{h^3},$$

i.e. the same expression as before. The improvement in the new method of counting, as compared with that of the Boltzmann statistics, is merely this, that the quantum theory assigns a definite size to the smallest cells, into which we have to partition the phase space for the purposes of statistics. In the ordinary kinetic theory of gases—disregarding the case of gas degeneration, which will be dealt with later—the size of the cells played no part, but dropped out in the further course of the investigation. In this case of light quantum statistics, however (as also in the case of gas degeneration, see § 5, p. 212), the size of the cells is of great importance. We may mention further, that the finite size  $h^3$  thus found for a cell corresponds precisely to Heisenberg's uncertainty principle, according to which the position and momentum of a particle cannot be more exactly defined than is consistent with the relation  $\Delta p \Delta q \sim h$ ; in view of this relation, it would in fact be quite meaningless to make a finer division of the phase space, as it is impossible to decide by experiment in which of these cells a particle lies.

We now return to the statistics of light quanta, and begin with a verification of the statement made above, that the idea of light quanta (together with the definition of the size of the cell by quantum considerations) is *not* sufficient to enable us to deduce Planck's formula, if we adhere to the Boltzmann statistics. For, according to Boltzmann's result, the number of quanta in a definite sheet—which we shall characterize by a mean  $\nu_s$  of the frequencies

of this sheet—is given by the product of the number  $g_s$  of the cells in this sheet and the Boltzmann factor  $Ae^{-\epsilon_s/kT}$ , so that

$$n_s = Ag_s e^{-h\nu_s/kT} \quad (\epsilon_s = h\nu_s).$$

The Boltzmann law of distribution was obtained, let us repeat (§ 6, p. 9), as the most probable distribution of the particles of a gas (in our case the light quantum gas) in the various sheets (called cells in our earlier investigation), subject to the two subsidiary conditions  $\Sigma n_s = n$  and  $\Sigma n_s \epsilon_s = E$  when the number of particles and the total energy are given. For the distribution of energy in our light quantum gas we therefore find

$$\begin{aligned} h\nu_s n_s &= Ag_s h\nu_s e^{-h\nu_s/kT}, \\ u_\nu d\nu &= \frac{h\nu n}{V} = A \frac{8\pi h\nu^3 d\nu}{c^3} e^{-h\nu/kT}, \end{aligned}$$

which, apart from the factor  $A$ , is Wien's law, instead of the result required by Planck's formula, viz.

$$u_\nu d\nu = \frac{8\pi h\nu^3 d\nu}{c^3} \frac{1}{e^{h\nu/kT} - 1}.$$

Thus we see that the classical statistical methods fail, not only in the case when we deal with the statistics of cavity radiation from the wave point of view (§ 3, p. 201), but also when we try to set up a statistical theory of the light quantum gas. The question therefore arises of what changes must be made in the classical statistics in order that it may become possible to deduce Planck's radiation law by purely statistical reasoning, without making use of the roundabout road by way of an absorbing and emitting oscillator.

We see in the first place that for  $T$  small, i.e. for  $e^{h\nu/kT} \gg 1$ , our formula would agree with Wien's approximation (valid in this case) to Planck's law, provided we took  $A = 1$ . What does the condition  $A = 1$  signify? If we go back to the proof of Boltzmann's law of distribution, we easily see that the value of the coefficient  $A$  follows from the first subsidiary condition (fixed number of particles). If this subsidiary condition were dropped, we would get  $A = 1$ ; the radiation formula which we have just deduced would then agree asymptotically with the correct formula. We must conclude therefore that the first condition does not hold for light quanta. A special reason justifying this conclusion is that for large quanta of energy, i.e. for large values of  $h\nu/kT$ , our method of proof (if only on account of the



correspondence principle) must give the correct result. Moreover, we can see at once that there is good reason for dropping the first subsidiary condition (fixed number of particles) in the case of light quanta; for of course in every emission process in an atom a new light quantum is formed, and in every absorption process one is absorbed and converted into other forms of energy.

But even this assumption is not sufficient by itself to allow Planck's formula to be established on a statistical foundation. We must in fact completely alter the statistical foundation itself, and set up a new "quantum statistics". The way in which this might be done was shown by Bose (1924) and improved by Einstein. They assume complete *indistinguishability* of the light quanta. The hypothesis is a very plausible one. Suppose the light quanta to be numbered in a certain order. If light quantum 1 is in the cell  $z_1$ , and light quantum 2 in the cell  $z_2$ , this distribution obviously represents the same state as the distribution in which light quantum 1 is in the cell  $z_2$ , and light quantum 2 in the cell  $z_1$ , while the distribution of the rest of the light quanta remains as before; for the two light quanta of course only differ from each other in just this, that one of them is in the first cell and the other in the second. Numbering or individualization of the separate light quanta is entirely meaningless, since the state is completely and uniquely described by merely specifying how many light quanta there are in cell  $z_1$ , how many in cell  $z_2$ , and so on. This means, however, an essential distinction as compared with the classical statistics. In the latter, two cases, either of which passes into the other by merely interchanging two light quanta, are counted in the enumeration of states as two different states, whereas now they represent the same state, and in the enumeration of states must be counted as only a single state.

We shall consider the subject from the mathematical standpoint, and that with just so much greater generality that in the first place we shall speak of particles in general, and not specially of light quanta alone. We have already seen in connexion with the statement of Pauli's principle (§ 5, p. 159) that there also it does not matter which precise electron is at this place or that; the second half of Pauli's principle was simply this, that the interchange of two electrons does not change the state of the system as a whole, so that two distributions, which only differ in respect of the interchange of two electrons, represent one and the same state.

From the standpoint of wave mechanics every particle, a light quantum included, is described by the specification of its proper func-

tion; let the proper function of one particle (the first) be  $\psi_k^{(1)}$ , that of a second  $\psi_l^{(2)}$ , of a third  $\psi_m^{(3)}$ , and so on; where  $k, l, m, \dots$  represent the state of the particle in question (in the case of light quanta, e.g.,  $k$  stands for the three quantum numbers  $k_1, k_2, k_3$ ). The state as a whole is then described (to a first approximation at least, neglecting mutual action between individual particles) by the product of these proper functions, i.e. by a proper function

$$\Psi_{klm\dots} = \psi_k^{(1)}\psi_l^{(2)}\psi_m^{(3)}\dots$$

If two particles are now interchanged, e.g. particle 1 and particle 2, another wave function is obtained,

$$\Psi'_{klm} = \psi_k^{(2)}\psi_l^{(1)}\psi_m^{(3)}\dots,$$

which obviously corresponds to the same value of the energy of the whole system, viz. to

$$\epsilon_{klm} = \epsilon_k + \epsilon_l + \epsilon_m + \dots$$

We obtain other wave functions, for the same value of the energy, by taking an arbitrary linear combination of those wave functions which arise from the one first written down by a permutation of the individual particles among themselves, i.e. by interchange of the arguments 1, 2, 3 of the separate functions:

$$\Psi_{klm\dots} = \sum_P a_{123\dots} \psi_k^{(1)}\psi_l^{(2)}\psi_m^{(3)}\dots,$$

where the sum is to be taken over all permutations  $P$  of these arguments, and the factors  $a_{123\dots}$  represent arbitrary constant coefficients.

In the sense of classical statistics these wave functions give as many different states as there are linearly independent wave functions among them. From the standpoint of the new statistics, however, those cases which arise from one another by mere permutation of particles belong to the same state. Hence the wave function which describes this state cannot change when the particles are permuted, or at most it can only change in its sign, since only quadratic forms, such as  $|\Psi|^2$ , are of any account so far as physical interpretation is concerned. Now it is easily seen that the only wave function of the form specified above, which does not change when the particles are permuted, is the one in which all the coefficients are equal to 1, i.e. the *symmetrical wave function*

$$\Psi_s = \sum_P \psi_k^{(1)}\psi_l^{(2)}\psi_m^{(3)}\dots$$

Another possibility, in which the sign of the wave function changes, but not the value of its square, is the *antisymmetrical form*

$$\Psi_a = \sum_p \pm \psi_k^{(1)} \psi_l^{(2)} \psi_m^{(3)} \dots,$$

where the  $+$  sign is to be taken for an even permutation of the particles, the  $-$  sign for an odd permutation. This antisymmetrical form is known from the theory of determinants—it is the expansion of the determinant

$$\Psi_a = \begin{vmatrix} \psi_k^{(1)} & \psi_k^{(2)} & \psi_k^{(3)} & \dots \\ \psi_l^{(1)} & \psi_l^{(2)} & \psi_l^{(3)} & \dots \\ \psi_m^{(1)} & \psi_m^{(2)} & \psi_m^{(3)} & \dots \\ \dots & \dots & \dots & \dots \end{vmatrix}.$$

No other functions exist which satisfy the requirement of indistinguishability.

We note further a special feature in the case of the antisymmetric function. We know that a determinant vanishes if two rows, or two columns, are the same; hence if two functions  $\psi_k$  and  $\psi_l$  are equal, the determinant vanishes, and with it the wave function of this state, i.e. this state does not exist. This is neither more nor less than Pauli's principle, that two electrons cannot be in the same state (i.e. cannot have their  $\psi_k$  the same).

There are therefore only two possible ways of describing a state by a wave function, viz. either by the symmetrical or the antisymmetrical wave function; the second possibility corresponds to Pauli's principle, the first is another, and entirely different matter. If we count the possible states on the basis of their wave functions (i.e. of the possible wave functions which are linearly independent), two different statistics present themselves. If we confine ourselves to the symmetrical wave functions (without Pauli's principle), we get the so-called *Bose-Einstein statistics*; if we describe the state by the antisymmetric function (with Pauli's principle), we get the *Fermi-Dirac statistics* (1926). Which of the two statistics we are to use in a particular case, it must be left to experience to decide. With regard to electrons, we already know that they obey Pauli's principle—we shall therefore deal with them tentatively by the Fermi-Dirac statistics (see § 6, p. 214); on the other hand, it turns out that we have to treat light quanta (Bose) and also gas molecules (Einstein) according to the Bose-Einstein statistics.

We proceed now to work out the latter type of statistics. In the first place we have to count the number of different states (i.e. of

linearly independent wave functions). In carrying out the enumeration, however, we use, not the wave picture, but the corpuscular picture, and have to find the number of distinguishable arrangements of the particles in a sheet, for the case of the *Bose-Einstein statistics*. For this purpose, we denote the individual cells of this sheet by  $z_1, z_2, \dots, z_{g_s}$ ; the number of them is, by definition, given by the weight factor  $g_s$  of this sheet. On the other hand, let there be  $n_s$  particles in this sheet, which for the present we denote individually by  $a_1, a_2, \dots, a_{n_s}$ . We have to distribute these particles among the  $g_s$  cells of the sheet, and determine the number of distinguishable arrangements. To this end, we describe a definite arrangement in the following way. We write down, purely formally, the elements  $z$  and  $a$  in an arbitrary order, e.g.

$$z_1 a_1 a_2 z_2 a_3 z_3 a_4 a_5 a_6 z_4 z_5 a_7 \dots,$$

with the understanding that the particles standing between two  $z$ 's are in each case supposed to be in the cell which stands to their left in the sequence; the sequence written down above means therefore that the particles  $a_1$  and  $a_2$  are in the cell  $z_1$ , the particle  $a_3$  in the cell  $z_2$ , the particles  $a_4, a_5, a_6$  in the cell  $z_3$ , no particle in  $z_4$ , and so on; that being so, the first letter in the symbolic arrangement must obviously be a  $z$ . We therefore obtain all possible arrangements, by first setting down a  $z$  at the head of the sequence—which can be done in  $g_s$  different ways—and then writing down the remaining  $g_s - 1 + n_s$  letters in arbitrary order one after the other. The total number of these arrangements is therefore

$$g_s(g_s + n_s - 1)!.$$

Distributions which can be derived from one another by mere permutation of the cells among themselves, or of the particles among themselves, do not, however, represent different states, but one and the same state; the number of these permutations is  $g_s! n_s!$ . We thus obtain for the number of distinguishable arrangements in the sheet which is characterized by the index  $s$ , in the case of the Bose-Einstein statistics,

$$\frac{g_s(g_s + n_s - 1)!}{g_s! n_s!} = \frac{(g_s + n_s - 1)!}{(g_s - 1)! n_s!}.$$

Altogether, the number of distinguishable arrangements for the case when there are  $n_1$  particles in the first sheet,  $n_2$  particles in the second, and so on, is given by the product of expressions of the above type, for the sheets:

$$W = \prod_s \frac{(g_s + n_s - 1)!}{(g_s - 1)! n_s!}.$$

We call this the "probability" of that distribution of the particles among the various sheets which is defined by the numbers  $n_1, n_2, \dots$ . It takes the place here of the probability found in the Boltzmann statistics (p. 12), viz.

$$W = \frac{n!}{n_1! n_2! \dots} g_1^{n_1} g_2^{n_2} \dots$$

The remaining part of the calculation proceeds as at the place cited. We have to determine the most probable distribution; for this purpose we use Stirling's theorem, and write

$$\log W = \sum_s \{ (g_s + n_s) \log (g_s + n_s) - g_s \log g_s - n_s \log n_s \},$$

where we have neglected the 1 in comparison with the large numbers  $g_s$  and  $n_s$ . We must now make  $\log W$  a maximum for variations of  $n_s$ , subject to the subsidiary condition

$$\sum_s n_s \epsilon_s = E, \quad (\epsilon_s = h\nu_s).$$

For light quanta, as we have shown above (p. 206), the second subsidiary condition (constancy of number of particles) drops out. Thus we find in the usual way

$$\frac{\partial \log W}{\partial n_s} = \log (g_s + n_s) + 1 - \log n_s - 1 = \log \frac{g_s + n_s}{n_s} = \beta \epsilon_s,$$

or

$$\frac{g_s + n_s}{n_s} = e^{\beta \epsilon_s}.$$

The Bose-Einstein law of distribution for light quanta therefore runs (if we drop the index  $s$ )

$$n = \frac{g}{e^{\beta \epsilon} - 1};$$

this gives for the energy density

$$u_\nu d\nu = \frac{n h \nu}{V} = \frac{1}{c^3} \frac{8 \pi h \nu^3 d\nu}{e^{\beta h \nu} - 1}.$$

This is just Planck's radiation formula, if we put  $\beta = 1/kT$ . The justification for this last step is given by thermodynamics; according to Boltzmann,  $S = k \log W$  is to be regarded as the entropy, and it can then be shown (see Appendix XXIX, p. 336) that from the equation  $T dS = dQ$  we can infer that  $\beta = 1/kT$  ( $dQ$  is the increment

of the heat content, or, at constant volume, of the energy content of the light quantum gas). From the Bose-Einstein statistics, therefore, Planck's radiation law can be deduced in a way to which no objection can be taken.

### 5. Einstein's Theory of Gas Degeneration.

After the brilliant success of the Bose-Einstein statistics with the light quantum gas, it was a natural suggestion to try it in the kinetic theory of gases also, as a substitute for the Boltzmann statistics. The investigation, which was undertaken by Einstein (1925), is based on the hypothesis that the molecules of a gas are, like light quanta, indistinguishable from each other.

The calculations run exactly as in the light quantum case, except that here a second subsidiary condition appears, on account of the conservation of the number of particles:

$$\sum_s n_s = N.$$

The determination of the probability of a definite distribution  $n_1, n_2, \dots$  follows the same lines as before. The calculation of the most probable distribution leads now, owing to the presence of the second subsidiary condition, to the equation

$$\frac{\partial \log W}{\partial n_s} = \log \frac{g_s + n_s}{n_s} = \alpha + \beta \epsilon_s,$$

or, on dropping the suffix  $s$ ,

$$n = \frac{g}{e^{\alpha + \beta \epsilon} - 1},$$

where again  $\beta = 1/kT$  (see Appendix XXIX, p. 336). Here the number  $g$  of cells in a sheet can be expressed by the corresponding energy; we have of course

$$\epsilon = \frac{1}{2m} p^2, \quad \text{and} \quad d\epsilon = \frac{1}{m} p dp,$$

where  $p$  is the momentum of the particles; the expression for  $g$  obtained above therefore becomes

$$g = \frac{4\pi V}{h^3} p^2 dp = \frac{4\pi V}{h^3} \sqrt{2m^3 \epsilon} d\epsilon.$$

We thus find the *Bose-Einstein law of distribution for atoms*:

$$\begin{aligned} dN &= F(\epsilon) \sqrt{\epsilon} d\epsilon \\ &= \frac{4\pi V}{h^3} \frac{\sqrt{2m^3} \sqrt{\epsilon} d\epsilon}{e^{\alpha + \beta\epsilon} - 1}, \quad (\beta = 1/kT), \end{aligned}$$

while the law of distribution given by the Boltzmann statistics was (p. 15)

$$\begin{aligned} dN &= Vdn = 4\pi Vn \left( \frac{m}{2\pi kT} \right)^{\frac{3}{2}} e^{-(m/2)(v^2/kT)} v^2 dv \\ &= 4\pi N \left( \frac{m}{2\pi kT} \right)^{\frac{3}{2}} \sqrt{\frac{2}{m^3}} e^{-\epsilon/kT} \sqrt{\epsilon} d\epsilon. \end{aligned}$$

( $N$  is the number of particles in the volume  $V$ ,  $n$  the number per unit volume.)

The quantity  $\alpha$  is of course determined from the subsidiary condition

$$\int dN = \int_0^\infty F(\epsilon) \sqrt{\epsilon} d\epsilon = N = nV.$$

The constant  $\alpha$ , or more usually  $A = e^{-\alpha}$  is called the *degeneracy parameter*, the reason for the name being as follows. If  $\alpha$  is very great, and therefore  $A$  very small (compared to 1), we can neglect the 1 in the denominator in comparison with  $e^{\alpha + \beta\epsilon} = e^{\beta\epsilon}/A$  (since of course  $\beta\epsilon$  is always positive); thus in this case we obtain

$$dN = (4\pi V/h^3) \sqrt{2m^3} \sqrt{\epsilon} d\epsilon A e^{-\beta\epsilon},$$

i.e. the classical distribution law (Maxwell); here  $A = e^{-\alpha}$  is determined at once from the subsidiary condition of constancy of the number of particles. We thus find (p. 15; Appendix I, p. 259)

$$A = \frac{n h^3}{(2\pi m k T)^{3/2}}.$$

If then  $A$  is very small compared to 1, the Bose-Einstein distribution formula passes over into the classical one. It is otherwise in the case when  $A$  becomes comparable with 1 (the case of  $A > 1$ , i.e.  $\alpha < 0$ , cannot occur, for then the denominator vanishes for the energy value  $\epsilon = -\alpha/\beta$ , and for smaller values of  $\epsilon$  becomes negative, so that the whole theory becomes meaningless); if  $A \sim 1$ , deviations from the classical properties occur; we say then that *the gas is degenerate*. In this case the subsidiary condition leads to a transcendental equation

for  $A$ , which can be solved by expanding in ascending powers of  $n\hbar^3/(2\pi mkT)^{3/2}$ , the expansion beginning with the expression given above for the limiting case  $A \ll 1$ :

$$A = \frac{n\hbar^3}{(2\pi mkT)^{3/2}} \left( 1 - \frac{3}{4} \frac{n\hbar^3}{(4\pi mkT)^{3/2}} + \dots \right).$$

By substituting particular values of the constants  $n$ ,  $m$  and  $T$  we can read off from this equation whether under these conditions the gas is degenerate or not. We see in the first place quite generally that  $A$  increases, and accordingly the degeneracy becomes greater as  $n$  increases, i.e. as the density increases; on the other hand, it diminishes with increasing temperature and atomic weight. To take a special example: for hydrogen gas under normal conditions (for  $T = 300^\circ \text{K.}$ ,  $n \sim 3 \cdot 10^{19} \text{ cm.}^{-3}$ ) approximately  $A \sim 3 \cdot 10^{-5} \ll 1$ ; for heavy gases  $A$  becomes still smaller; gases are therefore never degenerate at normal temperatures and pressures, but behave according to the classical laws. Degeneracy would only become noticeable at unattainably low temperatures or at extremely high pressures, that is to say, in regions in which even according to the classical statistics the gases no longer behave as ideal gases (influence of finite size of the particles, condensation of the gas, &c.). Thus the Bose-Einstein statistics, when applied to gases, in the region in which the kinetic theory of gases is valid, shows practically no differences as compared with the classical statistics of Boltzmann.

## 6. Fermi-Dirac Statistics.

We have shown in § 4 (p. 209) that the introduction of the principle of indistinguishability into statistics leads to two, and only two, new systems of statistics, one of which, the *Bose-Einstein* statistics, we have discussed in detail in the last two sections (light quanta, gas molecules). We turn now to the second possible statistics, which is based on Pauli's principle, and was introduced by Fermi and Dirac. We have seen in § 4 (p. 209) that this statistics is intimately connected with the employment of Pauli's principle, observing that the proper function of a state in which two electrons have the same partial proper function (with respect to the four quantum numbers, including the spin quantum number) automatically vanishes.

To deduce the nature of this statistics, we shall use the model of a gas consisting of electrons, which, on the experimental evidence of spectra, obey Pauli's principle. In this case again, our first aim is to find the distribution of the electrons over the individual cells, bearing in



mind, however, that there are now twice as many cells as in the previous case of gas atoms, in consequence of the two possible settings of the spin; on the other hand, no cell can be occupied by more than one electron, or, in other words, the "occupation numbers" of the cells in this case must be either 0 or 1. (We might of course proceed on the assumption that the number of cells in each sheet is only half as great, but in compensation for this provide two possible places for the electrons in each cell, corresponding to the two directions of spin.)

We begin as before with the *enumeration of the distinguishable distributions*. Let the number of electrons in the  $s$ th sheet be  $n_s$ , these being distributed over the  $g_s$  cells of this sheet; of these  $g_s$  cells,  $n_s$  are therefore (singly) occupied (1), and  $g_s - n_s$  are empty (0). We characterize such a distribution by assigning to each cell its occupation number:

$$\begin{array}{cccccccccccc} z_1 & z_2 & z_3 & z_4 & z_5 & z_6 & z_7 & z_8 & z_9 & z_{10} & \dots \\ 1 & 0 & 0 & 1 & 1 & 1 & 0 & 1 & 0 & 0 & \dots, \end{array}$$

or by specifying the cells which are occupied by no particle, and those occupied by 1:

$$\begin{array}{cccccc|cccc} & & & 0 & & & & & & & 1 \\ z_2 & z_3 & z_7 & z_9 & z_{10} & \dots & z_1 & z_4 & z_5 & z_6 & z_8 & \dots \end{array}$$

Clearly there are  $g_s!$  such distributions, corresponding to the permutations of the  $g_s$  cells  $z$  in this scheme. But the same state (as regards occupation) is denoted by all distributions among these, which only differ from one another by permutation of the  $n_s$  occupied cells, or of the  $g_s - n_s$  unoccupied cells. Hence the "probability" of a distribution characterized by the occupation numbers  $n_1, n_2, n_3 \dots$  for the individual cells is given by

$$W = \prod_s \frac{g_s!}{n_s! (g_s - n_s)!},$$

or, from Stirling's theorem, by

$$\log W = \sum_s \left\{ g_s \log g_s - n_s \log n_s - (g_s - n_s) \log (g_s - n_s) \right\}.$$

As before, we wish to find the most probable distribution subject to the two subsidiary conditions

$$\sum_s n_s = N, \quad \sum_s n_s \epsilon_s = E.$$

We obtain it in the usual way:

$$\begin{aligned}\frac{\partial \log W}{\partial n_s} &= -\log n_s + \log(g_s - n_s) \\ &= \log \frac{g_s - n_s}{n_s} = \alpha + \beta \epsilon_s,\end{aligned}$$

or

$$n_s = \frac{g_s}{e^{\alpha + \beta \epsilon_s} + 1},$$

i.e., except for the  $+$  sign in the denominator, the same formula as in the case of the Bose-Einstein statistics. This difference of algebraic sign, however, carries with it a fundamental distinction between the present case and that of the Bose-Einstein statistics, in that  $\alpha$  can now have all values from  $-\infty$  to  $+\infty$ , and the degeneracy parameter  $A = e^{-\alpha}$  all values therefore from 0 to  $+\infty$ ; for the denominator of the distribution function is here always greater than 1. If we now substitute the value of  $g_s$  (with the factor 2 on account of the two directions of spin), we find, in the same way as above (p. 212), the *Fermi-Dirac law of distribution*:

$$dN = F(\epsilon) \sqrt{\epsilon} d\epsilon = \frac{8\pi V}{h^3} \frac{\sqrt{2m^3} \sqrt{\epsilon} d\epsilon}{e^{\alpha + \beta \epsilon} + 1}, \quad (\beta = 1/kT).$$

The degeneracy parameter is determined as before from the first subsidiary condition:

$$\int dN = \int_0^\infty F(\epsilon) \sqrt{\epsilon} d\epsilon = N = nV.$$

This transcendental equation for  $A$  can be solved, in the case when  $A \ll 1$ , by means of a power expansion, as before (p. 214):

$$A = \frac{nh^3}{2} (2\pi mkT)^{-3/2} + \dots \quad (A \ll 1).$$

When  $A$  is extremely small, this case, it is easy to see, passes over into the classical statistics. In the case when  $A \gg 1$ , we find an approximate expansion of the form (which we merely state, without proof):

$$\log A = -\alpha = \frac{h^2}{2mkT} \left( \frac{3n}{8\pi} \right)^{2/3} + \dots \quad (A \gg 1);$$

this is the case of *degeneracy* of the electron gas. A more thorough discussion of these formulæ is given in the next section, where at the

same time the most important application of the Fermi-Dirac statistics will be dealt with.

## 7. Electron Theory of Metals. Energy Distribution.

In the preceding section we spoke of an "electron gas", and pictured it to ourselves as a definite number ( $n$  per cm.<sup>3</sup>) of electrons, moving freely, without mutual disturbance. Such a case is of course unrealizable, since in virtue of their electric charge the electrons will always act upon each other; however, to a first approximation we can neglect this disturbing action, owing to the neutralizing effect of the positive ions.

Now such an approximately *free electron gas* exists in the interior of a metal. In proof of this statement, we may cite the high conductivity of metals; and metallic conduction is carried on by electrons only, since, in contrast to electrolytic conduction, it is demonstrably not bound up with transport of matter. In order to secure that the electrons shall be able to react to an external electric field in such a way that the resulting conductivity shall be as high as it is, they must at least approximately be able to move freely in the metal, in contrast to what happens in non-conductors, in which the electrons are firmly bound to the atoms. Founding on these ideas, even the older theories of metallic conduction (Riecke, Drude, Lorentz, since 1900) were able to give a satisfactory explanation of the *Wiedemann-Franz law*, which states that the electrical and thermal conductivities are proportional to one another, and that their ratio is inversely proportional to the absolute temperature. The older theories, however, consistently led to difficulties in regard to the explanation of the specific heat of metals. It is established by experiment that metals obey the Dulong-Petit law, i.e. that their specific heat, referred to 1 mole, is 6 cal./degree. This could be explained at once, if the temperature of the metal were determined solely by the vibrational energy of the atoms in the lattice, since the mean energy per lattice point is  $3kT$ . But to explain the process of conduction and other related phenomena it is necessary to assume that to every atom (ion) there corresponds approximately one free electron. The free electrons take part in the thermal motion in the metal, and in fact (as shown by the Wiedemann-Franz law) are largely responsible for the high thermal conductivity of metals. According to classical statistics, every free electron in the metal would therefore possess the mean kinetic energy  $\frac{3}{2}kT$ , and the specific heat of the metal would then be, per atom, not  $3k$ , but of the order  $(3 + \frac{3}{2})k$ , i.e. referred to 1 mole, 9 cal./degree, which is contrary to the facts.

The solution of this difficulty is due to Pauli and Sommerfeld (1927), who pointed out that the laws of classical statistics ought not to be applied to the electron gas within a metal, since it is bound to behave as a *degenerate gas*. Thus, since the mass of the electron is 1840 times smaller than that of the hydrogen atom, it follows that, at room temperature ( $T = 300^\circ$ ) and an electronic density of  $n \sim 3 \cdot 10^{19}$ , corresponding to a gas density at a pressure of 1 atmosphere, the value of the degeneracy parameter  $A_e$  for the electron is

$$A_e = A_H^{1/2} (1840)^{3/2} = A_H 4 \cdot 10^4 \sim 1.2,$$

where  $A_H$  is the degeneracy parameter for hydrogen gas under the same conditions; even in this case, therefore,  $A$  is of the order of magnitude 1. Values decidedly greater than this are found for the electron gas in metals. In silver the number of atoms per cm.<sup>3</sup> is  $n = 5.9 \times 10^{22}$ . Since, as we have already remarked, we must assume that roughly speaking there is one free electron for each atom, we find for this value of  $n$  by the first approximate formula that  $A$  is about 2300; in this case, therefore, the gas is degenerate to a high degree. For a value of  $A$  so high as this, it is true, it is not permissible to apply the first approximate formula, and we must use the second; but even this gives the still high value  $A \sim 210$ . The electron gas in metals is therefore in all cases highly degenerate—its properties are essentially different from those of an ordinary gas.

The most important characteristics of the Fermi-Dirac distribution function are the slightness of the dependence of distribution on temperature, and the occurrence of a zero-point energy. The latter property is closely connected with the Pauli principle. In the classical gas theory the absolute zero is characterized by the fact that the mean kinetic energy of the gas particles vanishes at that temperature, and accordingly the energy of every individual particle also vanishes; classically, therefore, at the absolute zero the gas particles are at rest. It is different in the Fermi-Dirac statistics; here each cell can only have a single occupant; in the state of lowest energy, all the cells of small energy are occupied, and the limit of the "filling up" of the system of cells is given by the number of electrons. We characterize this limit by the momentum  $p_0$  of that cell, up to which the filling up reaches; it is found from the formula deduced above for the number of cells:

$$2 \frac{4\pi V}{3h^3} p_0^3 = N,$$

or 
$$p_0 = h \sqrt[3]{\left(\frac{3N}{8\pi V}\right)} = h \sqrt[3]{\frac{3n}{8\pi}}.$$

The limiting energy  $\epsilon_0$  is then given by

$$\begin{aligned}\epsilon_0 &= \frac{p_0^2}{2m} = \frac{h^2}{2m} \left(\frac{3n}{8\pi}\right)^{\frac{2}{3}} \\ &= 5.77 \times 10^{-27} n^{2/3} \text{ ergs} = 3.63 \times 10^{-15} n^{2/3} \text{ electron volts.}\end{aligned}$$

We therefore obtain the following distribution curve for the electrons at the absolute zero (see fig. 7). Taking the electronic energy  $\epsilon$  as abscissa, and as ordinate the previously (p. 216) defined distribution function  $F(\epsilon)$ , whose product by the factor  $\sqrt{\epsilon}d\epsilon$  gives the number of electrons with energy values between  $\epsilon$  and  $\epsilon + d\epsilon$ , we find for the graph a rectangle; up to the energy value  $\epsilon_0$  the cells are com-

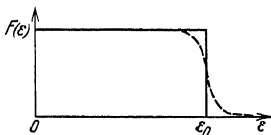


Fig. 7.—Fermi's distribution curve. The continuous, sharp-cornered line corresponds to the absolute zero ( $T = 0$ ), the dotted line to a temperature other than zero.

pletely filled, the cells with greater energy values are empty. In this case the degeneration parameter  $A$ , as the approximate formula shows, becomes infinitely great, as  $1/T$  in fact; comparison with the foregoing formula for the limiting energy shows that we can put approximately

$$\alpha = -\frac{h^2}{2mkT} \left(\frac{3n}{8\pi}\right)^{\frac{2}{3}} = -\frac{\epsilon_0}{kT}.$$

The distribution function, which is approximately valid for large values of  $A$ , i.e. for low temperatures, is then

$$F(\epsilon) = \frac{8\pi V}{h^3} \frac{\sqrt{2m^3}}{e^{(\epsilon - \epsilon_0)/kT} + 1},$$

and for the limiting case when  $T \rightarrow 0$ , gives the graph of fig. 7. When  $\epsilon < \epsilon_0$ , the exponential function in the denominator vanishes when  $T \rightarrow 0$ , and we have  $F(\epsilon) = 8\pi V \sqrt{2m^3}/h^3$ ; when  $\epsilon > \epsilon_0$ , however, the exponential function in the denominator becomes infinite, and  $F(\epsilon)$  vanishes.

As the temperature increases, the electrons are gradually raised into higher states; but the change in the electronic distribution will at first only take effect at the place where the Fermi function falls

away, and that by slowly rounding off the corners of the distribution curve, as indicated in fig. 7. The main body of the electrons, however, is left untouched by the rise of temperature. For not too high temperatures, therefore, only a vanishingly small fraction of the electrons takes part in the thermal motion, so that the specific heat of the electrons is very small. It is only when high temperatures are reached (of the order of  $10^4$  degrees C.), far above room temperature, that the tight packing of the electrons in the deeper states gradually becomes loosened, and we obtain a noticeable contribution from the electrons to the specific heat of metals.

### 8. Thermionic and Photoelectric Effect in Metals.

A further proof of the correctness of the idea of free electrons in metals and of the applicability of the Fermi-Dirac statistics is furnished by the phenomenon of thermionic

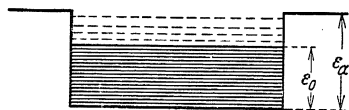


Fig. 8.—Diagram of the potential relations in a metal; the potential hollow is partly "filled up" with electrons.

emission. It is known that electrons emerge spontaneously from incandescent metals (e.g. glowing cathodes), and that these electrons in the absence of an applied potential form an electronic cloud or atmosphere round the incandescent body. Their number can be determined by measuring the current set up when an external E.M.F. is applied. Theoretically, the phenomenon of thermionic emission has to be pictured in the following way (fig. 8). Within the metal the electrons can certainly move freely, but in general their escape from the metal is opposed by a potential barrier, which is called the *work function*,  $\epsilon_a$ . At higher temperatures, however, it may happen that the energy of an electron becomes greater than  $\epsilon_a$ , so that it can escape from the metal. Using the formulæ of the Fermi-Dirac statistics, we can determine the number of electrons emerging in this way per unit time; the current is found to be

$$i = \frac{4\pi em}{h^3} (kT)^2 e^{-(\epsilon_a - \epsilon_0)/kT},$$

while the classical statistics gives for it the expression (Richardson, 1902)

$$i = en \sqrt{\frac{kT}{2\pi m}} e^{-\epsilon_a/kT},$$

which, regarded as a function of the temperature, differs somewhat from that given by the new theory (see Appendix XXX, p. 338). The two formulæ differ both in the power of  $T$  before the exponential function, and in the meaning of the constant in the exponent. To test them, it is usual to graph  $y = \log i$  against  $x = 1/T$ , i.e. the function

$$y = A - a \log x - bx \begin{cases} \text{Classical} & \text{Fermi Statistics} \\ A = \log(en\sqrt{k/2\pi m}) & \log(4\pi emk^2/h^3) \\ b = \epsilon_a/k & (\epsilon_a - \epsilon_0)/k \\ a = \frac{1}{2} & 2 \end{cases}$$

Now the term  $a \log x$  is in general so small compared to the other two terms, that it has been impossible up till now to decide whether the quantum theory formula with  $a = 2$  is superior to the classical one with  $a = \frac{1}{2}$ . On the other hand, the constant  $b$  is easily found, for, when the term  $a \log x$  is omitted, the equation  $y = A - bx$  represents a straight line, and  $b$  is its gradient. We can therefore determine the difference  $kb = \epsilon_a - \epsilon_0$  experimentally for various metals, and from this calculate  $\epsilon_a$  by substituting for  $\epsilon_0$  the value given by theory, viz. 0 for the classical, and  $3.63 \times 10^{-15}n^{2/3}$  electron volts for the quantum theory.

By comparing the values thus found for  $\epsilon_a$  with other measurements, it has been definitely settled that the quantum theory formula is correct, and not the classical one.

We can in fact determine the work function  $\epsilon_a$  in another way which is quite different, viz. by means of *diffraction experiments with slow electrons*. When a crystal lattice is irradiated with cathode rays, the electrons, as we know, are diffracted by the lattice, the position of the diffraction maxima depending upon the de Broglie wave-length  $\lambda = h/p$  of the electrons. When, however, the cathode ray enters the crystal, refraction occurs at the surface, since the kinetic energy of the electrons in the solid is greater, and therefore their wave-length smaller, than in the external space. As in optics, so here we can speak of a refractive index of the crystal with respect to electrons; this in fact, as in the optical case, being given by the ratio of the wave-lengths inside and outside; these are, if we reckon the energy from the bottom of the potential trough,

$$\lambda_{\text{outside}} = \frac{h}{p_a} = \frac{h}{\sqrt{2m(\epsilon - \epsilon_a)}},$$

$$\lambda_{\text{inside}} = \frac{h}{p_i} = \frac{h}{\sqrt{2m\epsilon}},$$

so that we find for the refractive index the expression

$$n = \frac{\lambda_{\text{outside}}}{\lambda_{\text{inside}}} = \sqrt{\frac{\epsilon}{\epsilon - \epsilon_a}}$$

By measurement of the position of diffraction maxima we can now determine the value of the refractive index and so, knowing  $\epsilon$ , calculate the work function  $\epsilon_a$ . For the sake of greater accuracy, we should choose  $\epsilon$  very small, i.e. use slow electrons.

Careful measurements of this kind by Davisson and Germer on single crystals of nickel gave  $\epsilon_a \approx 16$  electron volts. On the other hand, measurements on the thermionic effect give values of  $bk$  in the region of 4 electron volts, in disagreement with the classical formula, according to which we should have  $bk = \epsilon_a$ . The quantum theory, however, gives  $\epsilon_0 = 11.7$  electron volts, if we assume that in nickel two electrons per atom are free, in accordance with the fact that nickel has two valency electrons. This would give  $\epsilon_a - \epsilon_0 = 4.3$  electron volts, in good agreement with the results of measurement of thermionic emission. Measurements on a series of elements with respect to their refractive index have been carried out by Rupp, who found that we get good agreement between theory and experiment, if for the elements Cu, Ag, Au, Fe, Mo and Zr we assume two free electrons per atom, for Al 3, for Pb 3 or 4, while for K the number of free electrons per atom must be less than 2. But we cannot attach great weight to this agreement, seeing that we have made the rather rough approximation of treating the potential inside the metal as constant.

The same constants as occur in the theory of thermionic emission also determine the *photoelectric effect*, which sets in at the frequency given by  $h\nu = \epsilon_a - \epsilon_0$ , at which the energy of the incident light quantum is just sufficient to raise the electron from the limiting point of the Fermi distribution,  $\epsilon_0$ , to the height of the external potential. The same constants also determine the law of the *cold discharge*, in which the following state of affairs occurs. If by means of sharp points very high field strengths ( $10^6$  volts/cm.) are produced at the surface of a metal, electrons issue from the metal even though the temperature has not been raised. The explanation depends on the same principle as is applied in the theory of radioactive disintegration of the nucleus p. 182, and Appendix XXVI, p. 326). The external field implies a potential distribution in which the potential falls linearly from the surface outwards. We have therefore a potential barrier (fig. 9) at the surface of the metal, and we know that according to wave mechanics an electron can make its way through such a barrier. The greater



the external field, the narrower becomes the barrier, and therefore the greater becomes the number of electrons issuing per second. This number clearly depends on the height of the original potential above the zero-point level  $\epsilon_a - \epsilon_0$ , so that the number, if it could be found experimentally, would allow the value of  $\epsilon_a - \epsilon_0$  to be determined. The experiment, however, is impracticable, owing to the presence of minute inequalities and impurities in the surface, which change the magnitude of the field in an uncontrollable way; at every little projection the field is greater than the mean potential gradient. In point of fact, electrons begin to emerge at distinctly smaller field strengths than would be expected if the surface were ideally smooth.



Fig. 9.—Potential barrier at the surface of a metal.

### 9. Magnetism of the Electron Gas.

Another circumstance which confirms the correctness of our ideas about electrons in metals was pointed out by Pauli (1927). The electrons possess on account of their spin a magnetic moment equal to a Bohr magneton. As, according to our present picture, they behave in metals practically like free particles, one might expect, therefore, that they should cause a very strong paramagnetism. Experiment shows, however, that simple metals (e.g. Li, Na) are either not paramagnetic, or only very slightly so. Pauli explains this as follows. We can consider the valency electrons in the metal as free; the ions forming the residue, having an inert gas configuration of electrons, are non-magnetic. Of the free electrons (for  $T = 0$ ), two sit in every cell, and they have opposite spins, so that their magnetic moments exactly balance each other. If an external field  $H$  is applied, the electrons will tend to direct their spins parallel to the field, which they cannot do without leaving the doubly occupied cells and jumping into higher states. This increase in kinetic energy goes on until it is compensated by the decrease of potential energy due to the orientation in the field. As only a few electrons jump to higher states, the paramagnetism is much smaller than for systems not satisfying the exclusion principle of Pauli. When the temperature rises, the uppermost sheets of the Fermi distribution, as we have shown in § 7 (p. 219), begin to be "loosened", individual electrons being lifted out of the doubly occupied cells, so that now there are cells which are only singly occupied. But this gives, as can be shown,

only a second order effect. Consequently the paramagnetism of an electronic gas is almost independent of temperature. This is in agreement with experiment. But now the orbital motions of the electrons also produce a magnetic effect—diamagnetic indeed, as may be seen at once. The theory of this phenomenon is not simple, and there has been much discussion about it. From the standpoint of the quantum theory, however, a definite formula results for the magnitude of the diamagnetic moment. If we subtract this from the paramagnetic moment, we get the following table:

TABLE V—SUSCEPTIBILITIES OF THE ALKALI METALS

	Na	K	Rb	Cs
$\chi \cdot 10^7$ : Theory	4.38	3.40	3.26	3.02
$\chi \cdot 10^7$ : Experiment	5.8	5.1	0.6	-0.5

The agreement in order of magnitude is good. The differences, especially in Rb, and in Cs (where the sign does not agree), can be explained as due to the neglect of facts which should be considered. Thus, in the heavy elements the inner electronic shells make an appreciable contribution to the diamagnetism, and in all metals the assumption of entirely free electrons is incorrect, as we proceed to explain more fully.

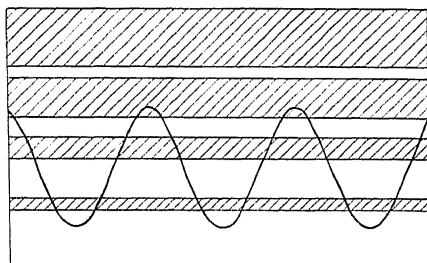
### 10. Electrical and Thermal Conductivity. Thermoelectricity.

To explain electrical conductivity, we must suppose the electrons in a metal to have a "free path". In fact, if we were to adhere in the problem of electrical conductivity to the idea of perfectly free electrons (that this is only a first approximation we have already emphasized above), the result would be an infinitely great conductivity. To explain finite resistance, therefore, we must take into account the fact that the electrons, in the course of their motion through the metal, collide from time to time with the ions of the lattice, and are thus deflected from their path, or are retarded; the average distance which an electron traverses between two collisions with the lattice ions is called, by analogy with the similar case in the kinetic theory of gases, the mean free path.

It has been shown by Sommerfeld (1928) that we can calculate the general behaviour of electrical and thermal conductivities without necessarily making special hypotheses as to the free path. The Wiedemann-Franz law follows from this theory; and we can explain in the same way the Joule heat, the Peltier and Thomson thermoelectric effects, and other phenomena.

The refinements of the theory, which have been worked out in particular by Houston, Bloch, Peierls, Nordheim, Fowler and Brillouin, have two main objects. In the first place, the picture of perfectly free electrons at a constant potential is certainly far too rough. There will be binding forces between the residual ions and the conduction electrons; we must elaborate the theory sufficiently to make it possible to deduce the number of electrons taking part in the process of conduction, and the change in this number with temperature, from the properties of the atoms of the substance. In principle this involves a very complicated problem in quantum mechanics, since an electron is not in this case bound to a definite atom, but to the totality of the atomic residues, which form a regular crystal lattice. The potential of these residues is a space-periodic function (fig. 10), and the problem comes to this—to solve Schrödinger's wave equation for a periodic potential field of this kind. That can be done by various approximate methods. One thing is clear: if an electron

Fig. 10.—The potential in a crystal lattice as a periodic function of position in space.



has little energy, and so lies deep in a hollow, the probability of its getting into another hollow is small; every hollow is isolated from the rest. This gives discrete energy levels. On the other hand, if an electron has a great deal of energy, and therefore is placed far above the level of all the hollows, it is practically free and has a continuous spectrum associated with it. Actually it will be a matter of energy strips which are narrow at the bottom and become gradually broader towards the top; finally we may speak of a continuum, with gaps always becoming narrower. We must remember, however, that the arrangement is three-dimensional; it may happen, therefore, that in one direction two neighbouring energy strips are completely separate, but overlap in another. Whether and when this happens, depends on the depth and arrangement of the potential hollows; if they are shallow, even the deepest energy strips will be broad, and in many directions run into one another.

These possible energy levels become filled up by electrons, two of which always fall into each state, on account of the spin. It turns out that such a strip contains exactly  $2N$  electronic states, if  $N$  is the

number of the electrons. If therefore one free electron is available per atom, as in the alkalis,  $N$  in all, the lowest strip is not yet completely occupied. The electrons will therefore be able to move freely, and we get a good conductor. In the alkaline earths, where there are two loosely bound electrons per atom, the first zone is completely filled. We then have an insulator in the case of deep energy troughs, where the next energy strip is completely separated from the lowest; but a conductor, when the troughs are so shallow that the first two strips overlap; for then both strips are practically empty, which means that the electrons can move freely. In this way all gradations of conductivity and its dependence on temperature can be explained qualitatively, and there exists to-day a satisfactory *theory of semi-conductors* (Fowler, Wilson).

We may mention one more phenomenon in metals, the *Hall effect*. When a current is flowing in a wire in which there is a transverse magnetic field, a potential difference occurs in the direction perpendicular to both the field and the wire. In itself this is not difficult to understand; for we know of course that a magnetic field exerts on a moving electron a force perpendicular to the field and to the velocity (Lorentz force, p. 26). What is remarkable is that in some metals this potential difference has its sign reversed, as if we were dealing with positive electrons. Positrons have now been discovered of course, but it cannot be these which are concerned here—in presence of so many negative electrons, their life is very short. But this anomalous Hall effect can be understood if we assume that one of the permitted energy strips is not completely filled up by (negative) electrons. The theory then shows that the “gaps” which remain behave exactly like free positive electrons.

We come now to the second main problem (cf. p. 224)—to determine by wave mechanics the free paths of the electrons—a problem which cannot be solved by the classical theory; it is a question of the scattering of the electronic waves which traverse the lattice of the metal, by the ions situated at the lattice points; and of the transference of their energy to the ionic lattice. The calculation gives thoroughly satisfactory results, bringing out correctly, for instance, the way in which electrical resistance depends on temperature. We may specially mention the calculation of the conductivity of alloys by Nordheim. It has been established experimentally that the resistance of an alloy is not always a monotonic function of the ratio in which the two components are mixed, but is generally greater for the mixture than for either of the pure metals; the investigation by wave

## VII] ELECTRICAL AND THERMAL CONDUCTIVITY

mechanics shows that in this matter an essential part is played the interference of the electronic waves.

In conclusion, let it be also mentioned, however, that up till now the theory is completely at a loss with regard to *supra-conductivity*, that is to say, the fact that many metals at very low temperatures—a few degrees above absolute zero—suddenly lose their resistance (Kamerlingh-Onnes, 1911); if, for example, an electric current is induced in a wire ring of supra-conducting material at a sufficiently low temperature, the current, in the absence of external disturbances, will continue with undiminished strength for a whole day. From the point of view of the theory, this phenomenon remains up to the present rather perplexing.

## CHAPTER VIII

### Molecular Structure

#### 1. Molecular Properties as an Expression of the Distribution of Charge in the Electronic Cloud.

This last chapter will be devoted to the subject of molecular structure. Our object in the first place will be to arrive at clear ideas with regard to two questions, viz. by what properties a molecule, from the physical standpoint, is most conveniently characterized and described, and how these properties can be determined experimentally. (Only after these points have been dealt with will we take up the question, how the phenomenon of chemical binding can be understood and explained physically.)

A molecule consists of a number of heavy nuclei, the atomic nuclei of the atoms or ions which form the molecule; round these nuclei the electrons revolve. Just as in the case of atoms, so here we can speak of an electron cloud. On account of the great difference between the masses of the electrons and the nuclei, the motion of the electrons is of course much more rapid than that of the nuclei—a circumstance which simplifies the discussion very decidedly. If, for example, we are investigating the motion of the electrons, we can to a first approximation regard the nuclei as at rest, since they move very little during the period of revolution of an electron. On the other hand, if we are examining the motion of the molecule as a whole (rotations), or of the individual atoms in the molecule relative to one another (vibrations), we can obtain a good approximation to the results by replacing all properties of the electronic motion by mean values; for during the time which must elapse before the nuclei have undergone any appreciable displacement from their original positions, the electrons make a great many complete revolutions.

With regard now to special molecular properties, one of the first importance is the *distribution of charge* in the molecule. As regards the total charge, we must distinguish here, as in the case of atoms, between neutral molecules and positive or negative ions. The charge distri-

bution itself is completely characterized by specifying on the one hand the mutual *distances of the nuclei*, and on the other the charge density  $\rho$  of the electrons. The latter can be regarded either classically as the mean charge per unit volume obtained by averaging over the motion of revolution of the electrons, or from the standpoint of wave mechanics as the charge density given by probability waves as in the case of the atom. The charge density on the one hand, in the case of equilibrium due to the action of the electrical forces, determines the nuclear distances; on the other hand, by its external boundary it gives the *molecular volume*  $v$  (from the standpoint of wave mechanics, in the case of a neutral molecule, exactly as in that of a neutral atom, the charge distribution falls off exponentially outwards, beyond a definite boundary, so that the size of the charge cloud can be assigned with comparative exactness).

We can now consider the question of the electrical centroids of the positive and the negative charges apart, that is to say, the electrical centroid of the nuclei alone, and the electrical centroid of the electron cloud alone. It may happen that the two points coincide, just as they always do in individual atoms, where the positive centroid is identical with the nucleus, and where also the centroid of the negative charge cloud, on account of the central symmetry of its charge distribution, always coincides with the nucleus. In general, however, the two centroids will be distinct from each other; consequently the external action of the molecule is like that of an electric dipole. In this case we speak of a *permanent electric dipole moment*, and denote it by the vector

$$\mathbf{p}_0 = \Sigma \bar{\mathbf{e}}\mathbf{r},$$

where the radius vectors  $\mathbf{r}$  are the position vectors of the nuclei and the electrons; the bar signifies averaging over the electronic motion, and the sum is taken over all the nuclei and electrons. If  $\mathbf{p}_0 = 0$ , the electrical behaviour of the molecule is determined to a first approximation by the *quadrupole moment*, defined by its components

$$\Theta_{xx} = \Sigma \bar{e}x^2, \dots, \Theta_{xy} = \Sigma \bar{e}xy, \dots;$$

this is a form exactly analogous to the mechanical moment of inertia; we therefore also speak of the *electrical moment of inertia*. It is a tensor, and as such can be represented by an ellipsoid (ellipsoid of inertia).

If the molecule comes into an electric field, it is deformed, the positive nuclei being attracted in the direction of the lines of force, the negative electrons in the opposite direction. Consequently, even

when there is no permanent dipole moment, a dipole moment is induced; to a first approximation its magnitude increases linearly with the field strength, and it can be represented by the relation

$$\mathbf{p} = \alpha \mathbf{E},$$

where  $\alpha$  is called the deformability or *polarizability*, and, as may easily be seen, has the dimensions of a volume. For spherically symmetrical molecules  $\alpha$  is of course a scalar, that is, a constant independent of direction. In general, however,  $\alpha$  depends on the direction, and may then be regarded as a tensor and represented by an ellipsoid, the *ellipsoid of polarization*. We may also mention the case where the molecules can rotate freely (gases); the ellipsoids of polarization corresponding to the individual molecules can then assume all possible positions in space, so that when an external field is applied, a mean polarization of the gas arises, which is given by  $\bar{\mathbf{p}} = \bar{\alpha} \mathbf{E}$ ; the mean value  $\bar{\alpha}$  over all directions is all that matters in this case.

As determining elements of molecules, therefore, besides the total charge  $e$ , the nuclear distances and the molecular volume, we have also now before us the dipole and quadrupole moments, as well as the polarizability. All these quantities (with the exception of  $e$ ), however, also depend more or less on the *state of excitation* of the molecule. Just as in the case of atoms, so also in molecules there exist different electronic states, characterized by quantum numbers which in the following sections we shall denote collectively by  $n$ . Moreover, the nuclear motion has an effect, consisting on the one hand of a rotation of the whole molecule, determined by the rotational quantum number  $j$ , and on the other hand of a vibration of the nuclei relative to one another, described by a vibrational quantum number  $s$ . With these different states of motion we shall deal in detail in the pages which follow.

## 2. Experimental Determination of the Molecular Constants.

We now proceed to explain in detail how the molecular constants enumerated in the preceding section can be determined experimentally. There is first the *molecular volume*, the determination of which, when neutral molecules are in question, can be effected by the methods of the kinetic theory of gases, already referred to in Chapter I (viscosity, free path, diffusion and direct measurement by molecular rays). The following table shows some molecular diameters \* so determined, in Ångströms:

\* The concept "molecule" in the kinetic theory of gases includes "monatomic molecule". Compare the concept "mole" (p. 3).



TABLE VI—DIAMETERS OF SOME MOLECULES, IN Å.

He 1.9	H <sub>2</sub> 2.3	H <sub>2</sub> O 2.6
Ne 2.3	O <sub>2</sub> 2.9	CO 3.2
A 2.8	N <sub>2</sub> 3.1	CO <sub>2</sub> 3.2
Kr 3.2	Cl <sub>2</sub> 3.6	C <sub>6</sub> H <sub>6</sub> 4.1
Xe 3.5		(C <sub>2</sub> H <sub>5</sub> ) <sub>2</sub> O 4.8

On account of their charge, the volume of ions has to be determined by other means. Two main methods have been used (Wasastjerna (1923), Goldschmidt (1926), Pauling (1927)). One deals with the "grating spaces" in ionic lattices, in rock salt, for example. If we assume that the molecules in crystal lattices are packed as tightly as possible, then the grating space gives directly the sum of the radii of the two ions, i.e. in the rock salt cube,  $r_{\text{Na}^+} + r_{\text{Cl}^-}$ ; for in the ionic lattices the ions are arranged in such a way that a positive ion is always surrounded by negative ions only, and conversely, so that the grating space is actually equal to the sum of the radii of the two ions. It is always only the sum of the two radii which we obtain in this way, not the radii themselves; if we knew one radius, we could then calculate all the rest. What we do is to measure grating spaces in crystals one of whose ions we have reason to believe is very small, for example,  $\text{Li}^+$ ; this has only two electrons, in the  $K$  shell, and will therefore be distinctly smaller than, for example, the  $\text{Cl}^-$  ion with completely filled  $K$  and  $L$  shells, and a full sub-shell (of 8 electrons) of the  $M$  shell. The grating space in the  $\text{Li}^+\text{Cl}^-$  lattice will therefore be approximately equal to the radius of the  $\text{Cl}^-$  ion.

The second method of determining ionic radii consists in measuring *ionic mobility* in electrolytes; small ions will make their way through the liquid more easily than large ones. The difficulty occurs in this method, however, that water molecules become deposited (hydration) on the ions, and so produce a deceptive appearance of substantially greater ionic radii. Here we give another table, showing ionic diameters of atomic ions, and for the sake of comparison we repeat the diameters of the inert gases; atoms or ions with similar electronic configurations are placed in the same row or column. We see that the negative ions, which have an inert gas configuration with a smaller nuclear charge than the corresponding inert gas, are larger than the latter, the reason being of course that the electrons in these ions are more loosely bound, so that their orbits have greater radii. A corresponding result, *mutatis mutandis*, holds for the positive ions also.

TABLE VII—DIAMETERS OF SOME ATOMIC IONS, IN Å.

H <sup>-</sup>	2.5	O <sup>--</sup>	2.6	S <sup>--</sup>	3.5	Se <sup>--</sup>	3.8
He	1.9	F <sup>-</sup>	2.7	Cl <sup>-</sup>	3.6	Br <sup>-</sup>	3.9
Li <sup>+</sup>	1.6	Ne	2.3	A	2.8	Kr	3.2
Be <sup>++</sup>	0.7	Na <sup>+</sup>	2.0	K <sup>+</sup>	2.7	Rb <sup>+</sup>	3.0
B <sup>+++</sup>	—	Mg <sup>++</sup>	1.6	Ca <sup>++</sup>	2.1	Sr <sup>++</sup>	2.5
C <sup>++++</sup>	0.4	Al <sup>+++</sup>	1.2	Sc <sup>+++</sup>	1.7	Y <sup>+++</sup>	2.1
N <sup>+++++</sup>	0.3	Si <sup>++++</sup>	0.8	Ti <sup>++++</sup>	1.3	Zr <sup>++++</sup>	1.7
		P <sup>+++++</sup>	0.7				

The second property to be considered is the *mean polarizability*  $\bar{\alpha}$ . Here we confine ourselves in the first instance to molecules without permanent dipole moment. By definition,  $\bar{\alpha}$  represents that mean dipole moment which is induced in a molecule by an electric field of unit strength (at least in the case of molecules which can rotate freely); the total polarization  $P$ , per unit volume (containing  $N$  molecules), in an external field  $E$ , is given by  $P = \bar{\alpha}NE$ . But, according to electrodynamics, the polarization  $P$  is connected with Maxwell's displacement vector  $D$  by the relation  $D = E + 4\pi P$ ; on the other hand, by definition,  $D = \epsilon E$ , where  $\epsilon$  is the *dielectric constant*. In the case of gases, where we can neglect the mutual action of the molecules, these relations lead to the following equation connecting the dielectric constant and the mean polarizability:

$$\epsilon = 1 + 4\pi N\bar{\alpha}.$$

In liquids, where the induced moments of the molecules influence each other, the relation is somewhat more complicated.

The dielectric constant can easily be measured by well-known methods, for instance by determining the refractive index  $n$  of the substance for long waves (infra-red), which by Maxwell's theory is, as we know, connected with  $\epsilon$  in this limiting case by the relation  $n \rightarrow \sqrt{\epsilon}$ . As has already been remarked, the whole line of argument is valid only for substances free from dipoles. We give a short table of mean polarizabilities for inert gases and atomic ions, forms of similar structure being placed in the same row or column. We can recognize here the same order of succession in the values as in the previous case of the diameters; this was to be expected, since of course to a greater diameter there corresponds a smaller binding force on the outer electrons, and accordingly a greater polarizability. We remark specially that  $\bar{\alpha}$  has the dimensions of a volume; it is always in fact of the same order of magnitude as the molecular volume.

TABLE VIII.—MEAN POLARIZABILITIES OF INERT GASES AND ATOMIC IONS. (The numbers denote  $\bar{\alpha} \cdot 10^{24} \text{ cm.}^3$ .)

He 0.202	F <sup>-</sup> 0.99	Cl <sup>-</sup> 3.05	Br <sup>-</sup> 4.17	I <sup>-</sup> 6.28
Li <sup>+</sup> 0.075	Ne 0.392	A 1.629	Kr 2.46	Xe 4.00
	Na <sup>+</sup> 0.21	K <sup>+</sup> 0.87	Rb <sup>+</sup> 1.81	Cs <sup>+</sup> 2.79
	Mg <sup>++</sup> 0.12		Sr <sup>++</sup> 1.42	
	Al <sup>+++</sup> 0.065			
	Si <sup>++++</sup> 0.043			

The polarizability of some neutral atoms (H, Li, K, Cs) has been recently determined by a method similar to the Stern-Gerlach experiment (§ 7, p. 166), namely, by measuring the deflexion of a beam of atoms in an inhomogeneous electric field (Stark, 1936). The results do not agree very well with theoretical computations from atomic models.

We go on now to consider molecules with *permanent dipole moment*  $\phi_0$ . Here, in addition to the polarization effect just considered, we have also the influence of the electric field on this permanent moment. In the absence of an external field, the moments of the individual molecules will have all possible directions, so that the gas is unpolarized. If an external field is applied, this tends to turn the individual dipoles round into the field direction (fig. 1); this ten-

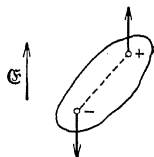


Fig. 1.—Couple exerted by an external electric field on a molecule with permanent dipole moment.

dency is opposed, however, by the thermal motion, which, as we have already frequently remarked, always has a smoothing out effect, and in this case acts in the sense of equal distribution of the dipole directions. Exactly the same relations are present here as in paramagnetism, where, as we have seen, it is a matter of the setting of the magnets in the direction of the magnetic field. For that case (see Appendix XXV, p. 322) we find a formula for the mean moment per unit volume, the field strengths being supposed not too great; we can apply that formula directly here, so that the polarization per unit volume due to the permanent dipoles is given by

$$P = \frac{Np_0^2}{3kT} E.$$

This is additional to the polarization determined by the polarizability of the molecules, and we therefore obtain a dielectric constant

$$\epsilon = 1 + 4\pi N \left( \bar{\alpha} + \frac{p_0^2}{3kT} \right) = \epsilon_0 + \frac{4\pi N p_0^2}{3kT},$$

where  $\epsilon_0$  stands for the dielectric constant for the case of a vanishing permanent dipole moment. The value of the dielectric constant therefore depends on two concurrent effects, one purely electrostatic and therefore independent of temperature, the other dynamical (orientation of the dipole moments) and a function of the temperature (Debye's law (1912), analogous to Curie's law for paramagnetism). By determining  $\epsilon$  exactly as a function of the temperature, we can therefore separate these two effects, and so from a series of measurements of the dielectric constant deduce the polarizability and the magnitude of the permanent dipole moment. Here, however, it is taken for granted that  $\epsilon$  is measured in the electrostatic way. For if we were to proceed as above by determining the refractive index (for infra-red waves), we should always be measuring the first effect only, that is, the one dependent on the polarizability. This is due to the fact that the orientation process in the dipole moments cannot follow the rapid vibrations of the electric vector of the light wave; for this orientation requires a rotation of the whole molecule and accordingly a motion of the atomic nuclei, which on account of the great mass of the nuclei takes place far too slowly to be appreciably affected by the rapidly varying electric forces of the light wave. Hence here also we have

$$n \rightarrow \sqrt{\epsilon_0} = \sqrt{1 + 4\pi N\bar{\alpha}},$$

so that a measurement of the refractive index always gives the polarization effect only. This makes it possible to determine  $p_0$  by a particularly simple method;  $n$  is measured optically and  $\epsilon_0$  found from it,  $\epsilon$  being then determined by a statical measurement at a known temperature (Debye):

$$\epsilon - \epsilon_0 = \frac{4\pi N p_0^2}{3kT}.$$

We note further that the limitation to long waves is necessary in the determination of the refractive index, in order to keep clear of the region of "anomalous dispersion", which is roughly characterized by the condition that the optical frequencies are of the same order of magnitude as the classical orbital frequencies of the electrons. In the long wave region, however, the electronic motion is much more rapid than the light vibration, so that the action of the light on the electrons depends only on the mean distribution of the electronic charge, or the polarizability.

The results so obtained have recently been successfully verified for some substances by direct deflection of a molecular beam (§ 7,

p. 17) in a non-homogeneous electric field (Estermann, 1928).

The following table contains a small selection from the very large number of published determinations of moments:

TABLE IX—DIPOLE MOMENTS OF MOLECULES (in  $10^{18}$  e.s.u.)

Carbon monoxide	..	CO	0.12
Carbon dioxide	..	CO <sub>2</sub>	0.0
Water	..	H <sub>2</sub> O	1.8
Methane	..	CH <sub>4</sub>	0.0
Methyl chloride	..	CH <sub>3</sub> Cl	1.9
Methylene chloride	..	CH <sub>2</sub> Cl <sub>2</sub>	1.6
Chloroform	..	CHCl <sub>3</sub>	1.0
Carbon tetrachloride	..	CCl <sub>4</sub>	0.0

With reference to this table we add some remarks. In general, molecules of symmetrical structure, such as CO<sub>2</sub>, CH<sub>4</sub>, CCl<sub>4</sub>, &c., have no dipole moment. Since H<sub>2</sub>O possesses a dipole moment, its structure cannot be symmetrical. At the present time it is assumed to have the form of an isosceles triangle (see § 6, p. 255). CO, as an asymmetric molecule, has of course a dipole moment. The series of chlorine compounds between methane and carbon tetrachloride is interesting;

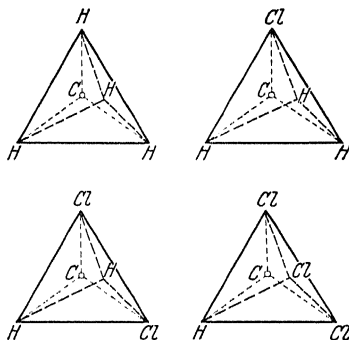


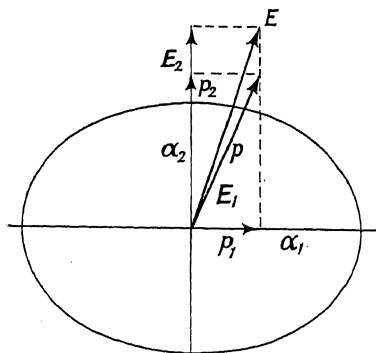
Fig. 2.—Structural formulæ of CH<sub>4</sub>, CH<sub>3</sub>Cl, CH<sub>2</sub>Cl<sub>2</sub>, and CHCl<sub>3</sub>. Methane is symmetric, and therefore has no dipole moment; the other three compounds are asymmetric, and more or less electrically polar.

their structural formulæ are shown in fig. 2. We see that the highly symmetric forms CH<sub>4</sub> and CCl<sub>4</sub> at the ends of the series have no dipole; in the intermediate cases dipoles are present, the gradations of which can be represented roughly in terms of "vectorial composition" of elementary dipoles.

We have already mentioned above that in the case of a gas free from dipoles ( $p_0 = 0$ ) the electric field arising from a molecule, and consequently the interaction between the molecules, are determined by the *quadrupole moment* (electrical moment of inertia). This therefore makes a contribution in this case to the *cohesive forces* which act between the molecules, and which in gases find their expression in the constants of the equation of state (e.g. van der Waals' equation, p. 20); they can be measured either by means of these constants or

from the latent heat of vaporization, &c. We do not consider this further here, however, but merely refer to § 7, p. 255, where we shall return to the matter.

We must add a remark with regard to polarizability. In what precedes we have taken account only of the mean value of  $\alpha$  over all directions—a procedure which, in the case of a gas, whose molecules can rotate freely, is certainly permissible as a first approximation. But by suitable experiments we can also determine the anisotropy of the polarizability, and so also form for ourselves a picture of the *anisotropy of the electron cloud*. We have already mentioned (p. 230) that the polarizability is a tensor, and can be represented by the so-called ellipsoid of polarization (see fig. 3). This has the following



properties. The three principal axes of the ellipsoid,  $\alpha_1$ ,  $\alpha_2$ ,  $\alpha_3$ , lie in the directions of the least and greatest polarizability, and the direction perpendicular to both of these; if electric field strengths of unit amount act in succession in these three directions, the lengths of the axes give the electrical

Fig. 3.—Diagram of the relative position of field and dipole moment in the case of an anisotropic molecule.

dipole moments corresponding to the respective cases. If we now let the unit electric field strength act obliquely to these three special directions, we can determine the polarization of the molecule by splitting up this field strength into its components in the three special directions, and determining the polarization effects of these three components separately; the total moment is then found by vectorial addition, and it is clear that in general the direction of the acting field strength does not agree with the direction of the induced dipole moment.

This has a marked effect on the polarization relations in the case of the *scattering of light*. Let us consider in the first place (see fig. 4) the case of the isotropic molecule or atom (atoms are in this sense always isotropic). If a light wave falls on this molecule, the electric vector  $E$  of the light wave excites in the molecule an electric moment  $p$ , which is parallel to the exciting field strength, and in phase with it; a scattered wave is therefore emitted having the same frequency as

the primary light. If we observe the scattered light in a direction perpendicular to the incident beam, we find that it is completely polarized; this is easy to understand, for its electric vector, which is determined by  $\mathcal{P}$ , is always parallel to the vector  $\mathbf{E}$  of the primary beam. This is not so for the anisotropic molecule, in which, as we have just seen, the induced dipole moment in general has a different direction from the exciting electric vector. If in this case again we observe the scattered light in a direction at right angles to the incident beam, we find that it is no longer completely polarized, but only partially; in the scattered light there now occurs a component of the electric vector at right angles to the incident light vector  $\mathbf{E}$ . If we make an experiment with the polarization apparatus, we no longer find, as we did before, a position of the nicol for which the field of view is completely dark. In this case we speak of the *depolarization* of the scattered light (Born (1917), Cabannes, Gans). By measuring the degree of depolarization, we can draw conclusions with regard to the anisotropy of the polarizability; for example, in the case of axisymmetric molecules (in which two of

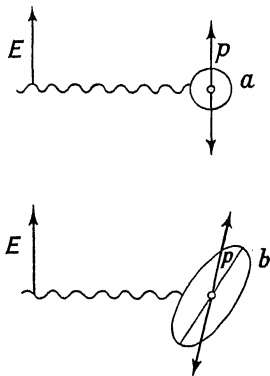


Fig. 4.—Depolarization due to the anisotropy of the polarizability. (a) In an isotropic molecule the induced dipole moment vibrates in the direction of the electric vector of the light wave. (b) In an anisotropic molecule it vibrates obliquely to  $\mathbf{E}$ , and causes depolarization in the scattered light.

the axes of the ellipsoid of polarization become equal:  $a_1 = a_2$ ), we can find, as an exact analysis shows, the value of  $a_3 - a_1$ . Since we can determine the mean polarizability by other measurements, we obtain in this way complete information regarding the lengths of the axes of the ellipsoid of polarization.

The same result can be obtained from the *Kerr effect* (1875). In a statical field  $\mathbf{E}$ , even in the absence of a permanent dipole moment ( $p_0 = 0$ ), anisotropic molecules are subjected to a couple, since the direction of the induced dipole moment does not fall in the field direction. This couple is again in the present case opposed by the thermal motion, which tends to produce a uniform distribution of direction, so that the molecules become partially oriented in the field direction, to an extent dependent on the temperature (Langevin, 1905; Born, 1916). But it can be shown that such a substance behaves towards light passing through it exactly like a doubly refracting

uniaxial crystal. In this case, again, measurement of the double refraction gives for axisymmetric molecules the value of  $\alpha_3 - \alpha_1$  (Gans (1921), Cabannes, Raman, Stuart). We may remark further that the effect in question is much employed in modern technical work; it is the basis of the Kerr cell (Carolus), extensively used in television technique as a light relay.

We come now to the determination of the distances between the nuclei, the frequencies of the nuclear vibrations, and other molecular properties connected with the nucleus. Here optical methods of a special sort play a part which we shall consider in next section.

### 3. Band Spectra and the Raman Effect.

We disregard in the first place the relative motions of the nuclei. A diatomic molecule, so far as its mass distribution is concerned, can be pictured as a nearly rigid dumb-bell, since of course the electrons by reason of their vanishingly small mass form an inappreciable factor in the mass distribution. This dumb-bell can turn round an axis fixed in space, and so possesses angular momentum, which according to Bohr must be quantised. If  $j$  is the quantum number of this angular momentum, the energy of the rotating dumb-bell is given on Bohr's theory (p. 100) by

$$E_j = \frac{\hbar^2}{8\pi^2 A} j^2 \quad (j = 0, 1, 2, \dots)$$

or, according to quantum mechanics (p. 129)

$$E_j = \frac{\hbar^2}{8\pi^2 A} j(j+1).$$

We have called the latter energy term the Deslandres term (as contrasted with the Balmer term). Here  $A$  is the moment of inertia of the dumb-bell about an axis through the centroid at right angles to the line joining the nuclei, and is easily found in terms of the nuclear distance and the masses of the two atoms. Thus, if  $r_1, r_2$  are the distances from the centroid of the atoms of masses  $m_1, m_2$ , the moment of inertia is by definition  $A = m_1 r_1^2 + m_2 r_2^2$ ; also  $m_1 r_1 = m_2 r_2$  and  $r_1 + r_2 = r$ , where  $r$  is the nuclear distance, which we wish to find. If then we denote by  $m$  the effective mass, i.e. if we take

$$\frac{1}{m} = \frac{1}{m_1} + \frac{1}{m_2}, \quad \text{or} \quad m = \frac{m_1 m_2}{m_1 + m_2},$$



we have for the moment of inertia

$$A = mr^2.$$

The nuclear distance can therefore be found if we know the values of the Deslandres terms, which can be determined from the emitted *rotational band spectra*. We have already pointed out (p. 101) that the spectrum emitted by a rotator consists of a series of equidistant lines. In fact, as we have already repeatedly stated, in simply periodic motions there is a selection rule  $\Delta j = \pm 1$ , so that the emitted frequency is found as the difference of two consecutive energy terms:

$$\begin{aligned} \nu &= \frac{E_j - E_{j-1}}{h} = \frac{h}{8\pi^2 A} \{j(j+1) - (j-1)j\} \\ &= \frac{h}{4\pi^2 A} j. \end{aligned}$$

By measuring the separation of the lines,  $\frac{h}{4\pi^2 A}$ , we therefore get  $A$ , and accordingly  $r$  also. It is assumed, however, in this method of deduction that the initial and final states of the electronic system are identical, since a difference in these would involve a change in the nuclear distance also; moreover, we have disregarded any possible oscillations of the nuclei relative to one another. The *purely rotational bands* are not very suitable in practice, however, for the determination of nuclear distances, since they lie in the extreme infra-red. We can easily make a rough estimate of their position. Thus, atomic masses are of the order of magnitude  $10^{-22}$  to  $10^{-23}$  gm., the nuclear distances are about  $10^{-8}$  cm., giving moments of inertia of approximately  $10^{-37}$  to  $10^{-39}$  gm. cm.<sup>2</sup>. We thus find for the frequencies values from  $10^9$  to  $10^{11}$  sec.<sup>-1</sup>, and therefore fractions of a centimetre for the wave-lengths.

TABLE X  
NUCLEAR DISTANCES AND MOMENTS OF  
INERTIA OF THE HYDROGEN HALIDES

	$r, 10^8$ cm.	$A, 10^{40}$ gm. cm. <sup>2</sup>
HF	0.93	1.35
HCl	1.28	2.66
HBr	1.42	3.31
HI	1.62	4.31

Above we give a table of a few nuclear distances and moments of inertia found by means of the infra-red bands.

In polyatomic molecules, we have different moments of inertia about different axes, and must determine them separately.

The circumstances become distinctly more favourable, if we also take into account the *nuclear vibrations*. We spoke at the outset of equilibrium of the forces between the nuclei and the mean distribution of the charges of the electrons. Round the position of equilibrium, which of course must be stable, the nuclei can oscillate, and the whole electron cloud pulsates along with it. The whole molecular energy, after deduction of the kinetic energy of the nuclei, is to be regarded as *potential energy*,  $V(r)$ , of the nuclear motion; this therefore includes, besides the pure Coulomb energy of the (positively charged) nuclei, the mean electronic energy, or, more exactly, the averaged energy of the electronic motion, calculated on the supposition that the nuclei are kept fixed. Rotations of the molecule as a whole are in the first instance disregarded. The equilibrium position of the nuclei is defined by the minimum of  $V(r)$ ; hence the equilibrium nuclear distance  $r_0$  is given by  $(dV/dr)_{r_0} = 0$ . Such a minimum necessarily exists, as otherwise no molecule having a finite nuclear distance could be formed at all. Fig. 5

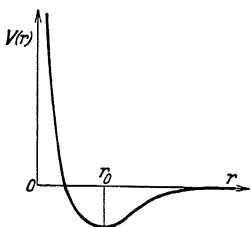


Fig. 5.—Graph of potential as a function of the distance between the two atoms combined in a molecule. The position of equilibrium is at  $r_0$ .

shows diagrammatically a potential curve of the type in question. From the minimum at  $r_0$  the potential rises very steeply as  $r$  becomes smaller, the Coulomb repulsion between the two nuclei preponderating here; in the direction of greater nuclear distances the potential curve flattens and asymptotically approaches a definite limiting level, which in the diagram has been arbitrarily taken as the zero level; this corresponds to the case of nuclei which are far apart, so that the molecule is practically completely split up into its constituent parts. In this position the potential energy  $V(r)$  of the nuclear motion is simply equal to the constant electronic energy of the two separated atoms, and can therefore be normalized to zero; the nuclei (and with them the whole atoms) then move as free particles.

In the potential hollow of  $V(r)$  the nuclei can vibrate as *quantum oscillators*. In the neighbourhood of the equilibrium position  $r_0$ , the potential curve is approximately parabolic in form; this we see from the Taylor expansion of  $V(r)$  at the point  $r = r_0$ :

$$V(r) = V(r_0) + \frac{(r - r_0)^2}{2} \left( \frac{d^2 V}{dr^2} \right)_{r_0} + \dots$$

For not too great amplitudes, therefore, the nuclei vibrate like harmonic oscillators, since of course in this case the restoring force is proportional to the distance; we can therefore apply the formulæ for the harmonic oscillator, which from the classical standpoint lead to the energy levels

$$E_s = h\nu_0 s \quad (s = 0, 1, 2, \dots),$$

and from the standpoint of wave mechanics to the formula

$$E_s = h\nu_0 (s + \tfrac{1}{2});$$

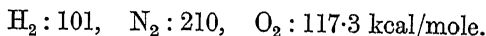
the first formula was deduced in § 2 (p. 101), the second is obtained by solving the wave equation of the harmonic oscillator (p. 124, or Appendix XVI, p. 295). The proper frequency  $\nu_0$  is determined by the restoring force, being given, as we may easily prove, by the equation

$$\nu_0 = \frac{1}{2\pi} \sqrt{\frac{1}{m} \left( \frac{d^2 V}{dr^2} \right)_{r_0}}.$$

By the rules of differential geometry  $(d^2 V/dr^2)_{r_0}$  is equal to the curvature of the potential curve at the point  $r = r_0$ , so that the result can also be expressed in the following form: the greater the curvature of the potential curve of the nuclear motion at the position of equilibrium, the greater is the proper frequency, and the higher are the corresponding energy levels.

The foregoing formulæ, as has already been mentioned, only hold for small amplitudes of vibration, or, what comes to the same thing, for the low quantum numbers. In the case of the more highly excited states, the deviation of the potential curve from the parabolic form has the effect of making it no longer allowable to treat the vibration as that of a harmonic oscillator; the formulæ deduced above have to be supplemented by corrections, which alter the values of the terms; when the quantum number  $s$  increases, the terms in fact crowd more and more closely together as they approach the so-called *convergence limit* (fig. 6). This limit corresponds to the *dissociation* of the molecule; it requires a quantity of energy equal to the depth of the potential hollow below the asymptotic limiting value of  $V(r)$ , that is, the quantity  $V(\infty) - V(r_0)$ ; on excitation by this or a greater quantity of energy the molecule splits up into atoms or ions, which then move apart with

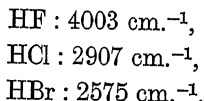
a definite velocity, given by the excess energy. In the spectrum, this finds its expression in the fact that a band occurs with a convergence limit, which is immediately followed by a continuum (Franck). From the position of the limit we can determine the *dissociation energy*, and that much more exactly than by the thermal measurements used in chemistry. Thanks to this principle, the work of separation is now well known for a great number of molecules, for example,



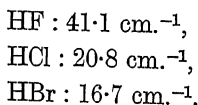
To give some idea of the order of magnitude of the frequencies of vibration, we quote here a few funda-

Fig. 6.—Term scheme of a vibrational band. The discrete energy levels converge towards a limit, which corresponds to the dissociation of the molecule; immediately above this a continuum follows, signifying that after dissociation the components of the molecule have kinetic energy, and fly apart.

mental vibrational frequencies  $\nu_{\text{osc.}}$  (after Czerny), in the form of wave numbers:



For comparison, we give also the corresponding fundamental rotational frequencies  $\nu_{\text{rot.}}$ :



It will be observed that the vibrational quanta are very decidedly larger than the rotational quanta.

Up till now, in dealing with vibrations we have disregarded the possible occurrence of molecular rotations. When absorption takes place, rotation and vibration can of course be excited simultaneously; the energy is then given approximately by

$$E_{s,j} = E_s + E_j = h\nu_0(s + \frac{1}{2}) + \frac{h^2}{8\pi^2 A_s} j(j+1);$$

the energy quanta of the rotational motion are for a first approximation simply added to the vibrational energy terms. The frequencies, which arise from differences of such terms (the selection rules  $\Delta j = 0, \pm 1$ ;  $\Delta s = \pm 1, \pm 2, \dots$  hold here), that is, for instance, for the fundamental vibration

$$\nu = \frac{E_{s,j} - E_{s-1,j-1}}{h} = h\nu_0 + \frac{h^2}{8\pi^2} \left\{ \frac{j(j+1)}{A_s} - \frac{j(j-1)}{A_{s-1}} \right\},$$

are the frequencies which give the rotational-vibrational bands; in consequence of the large value of the proper frequency  $\nu_0$  they lie in the short wave infra-red. The addition of the vibrational quantum  $h\nu_0$  therefore causes the whole spectrum to be displaced in the direction towards shorter waves, i.e. into the region which can be reached fairly easily by experimental methods; still, the law of succession of the rotational lines is more complicated here than in the case of pure rotations, since the moment of inertia, which is determinative for the separation of the lines, depends on the value of the nuclear distance for the time being, and may be different in the initial and final states.

The energy terms  $E_{s,j}$  written down above do not give the complete scheme of terms, since we laid down the condition that the electronic state of the molecule does not change in the rotational and vibrational transitions. But the electronic state can be altered by absorption or emission of light; for the molecule, exactly as for the atom, there are excited states, which are distinguished by quantum numbers  $n = 1, 2, \dots$ . To every such state there corresponds a particular potential energy of the nuclear motion,  $V_n(r)$ , which is found by averaging; we therefore obtain different curves for the potential of the nuclei, corresponding to the separate excitation levels of the electronic motion. The horizontal asymptotes ( $r \rightarrow \infty$ ) of these curves give the differences of energy in the end products occurring in dissociation. If, for example, molecular hydrogen  $H_2$  dissociates, it splits up, according to the electronic state of the molecule, either into two hydrogen atoms in the ground state, or into an unexcited and an excited hydrogen atom, or into two excited atoms; to the various energies of the products of dissociation

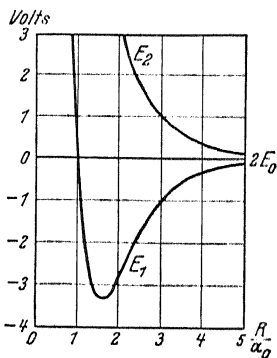


Fig. 7.—Potential curves for homopolar binding ( $H_2$  as example). One curve has a potential minimum (attraction); the other corresponds to a pure repulsion.

there correspond the various horizontal asymptotes of the potential curves for the nuclear vibrations.

In fig. 7 two potential curves are shown, one of which,  $E_1$ , possesses a minimum, and therefore makes a stable chemical binding possible, while the other,  $E_2$ , steadily falls; the latter of course does not corre-

spond to any chemical binding, but to a repulsion between the atoms, since its lowest state represents the state of atoms infinitely far separated. This double possibility occurs even in the case of the hydrogen molecule (see § 6, p. 253).

In the case of simultaneous excitation of higher electronic states, of vibration and of rotation, the total energy is given approximately by the formula

$$E = E_n + E_s + E_j;$$

where  $E_n$  represents the pure electronic energy, as determined by the energy difference between the minima of the various potential curves. The presence of  $E_n$  has the effect of displacing the bands into the visible or ultra-violet region, since the order of magnitude of the frequency determined by transitions from one electronic state of the molecule into another is the same as that of electronic transitions in the case of atoms; the combination of an electronic jump with a transition in the vibrational and rotational state implies the emission of a rotational-vibrational band in the region of wave-lengths fixed by the electronic transition. The appearance of these bands is the same as that of the pure rotational-vibrational bands, except for the fact that small "perturbations" occur, which are due to the interaction between the electronic and the nuclear motion (for example, alteration of the proper frequency of the nuclear vibration by an electronic jump) (fig. 8, Plate VIII, p. 168).

The electronic terms  $E_n$  for diatomic molecules admit of similar classification to those of atoms. Here, however, we cannot use the orbital angular momentum for this purpose, as in the case of atoms ( $l = 0, 1, 2, \dots$ ;  $S, P, D \dots$  terms), since the electron cloud has no longer a fixed total angular momentum, as it has in atoms; for the line joining the nuclei represents a specially distinguished direction in the molecule, and it rotates in space (rotational terms), carrying of course the electron cloud along with it. The *component of the electronic angular momentum* in this special direction is quantised, however, not being affected by the rotation of the molecule as a whole. The quantum number of this angular momentum component is denoted by  $\lambda$ , and, in analogy with the atomic case, the terms are designated by Greek capital letters, corresponding to the value of the quantum number:

$$\begin{array}{ccccccc} \lambda = & 0 & 1 & 2 & \dots & & \\ & \Sigma & \Pi & \Delta & \dots & \text{terms.} & \end{array}$$

We may also mention that special features appear in the scheme as

developed so far, if a molecule consists of two equal nuclei; degeneracies then occur, which express themselves in definite typical alterations in the spectrum (dropping out of certain lines).

Finally, we must refer to a complication which arises from the fact that the nuclei are really not point charges, but have a structure. The most important effect of this is the nuclear spin. The inner angular momenta of the nuclei have to be added vectorially to the angular momentum of the rest of the system, due to the rotational motion of the nuclei and electrons, with their spin. In the case of two equal nuclei, very large degeneracy effects arise from this cause. The simplest case is that of the molecule  $H_2$ . Here we have as nuclei two protons, each with the spin  $\frac{1}{2}$ . Vectorial composition gives  $\frac{1}{2} - \frac{1}{2} = 0$ , or  $\frac{1}{2} + \frac{1}{2} = 1$ ; in the latter case the resultant nuclear moment 1 can have three settings with respect to the moment of the rest of the motions, determined by the possible components  $-1, 0, +1$  of the vector of length 1. These states have all the same probability. Hence, molecules with nuclear moment 1 will be three times more frequent than those with the nuclear moment 0. Further, it is found that practically no transitions take place spontaneously from the one sort of molecule to the other (Heisenberg, Hund); they exist almost independently of each other, and so they have been given names. Molecules with the moment 0 are called parahydrogen; those with the moment 1, orthohydrogen. Ordinary hydrogen is a mixture of these in the ratio 1 : 3, as is shown by many properties, especially the specific heat (Dennison, 1927); but orthohydrogen, which has the higher energy content, can be converted into parahydrogen by means of catalytic effects, such as adsorption at surfaces (Bonhoeffer and Harteck, Eucken and Hiller, 1929). The difference in respect of energy arises from the fact that the paramolecules with the spin 0 can only have even rotational quantum numbers, and orthomolecules with the spin 1 only odd ones; the lowest state with quantum number 0 belongs therefore to parahydrogen. Since the latter occurs in ordinary hydrogen in the ratio 1 : 3 as compared with the ortho-modification, the intensities of the band lines are in the same ratio. Such a variation in the intensity of band lines occurs in all molecules which consist of two equal atoms with nuclear spin, and makes it possible to determine this spin. The method has already been mentioned (p. 178).

Band research has already developed into a science of considerable dimensions. The investigation and analysis of band spectra gives very far-reaching information regarding the structure of molecules. For the interpretation of the extensive experimental material which has

been collected, ingenious mathematical methods, such as group theory, considerations of symmetry, and so on, have been called into service.

To conclude this section, there is one other phenomenon we should like to discuss, viz. the *Raman effect*. Let it be mentioned beforehand, however, that this is not a revolutionary discovery, like, for example, the discovery of the wave nature of the electron, but an effect which was predicted by the quantum theory (Smekal (1923), Kramers-Heisenberg) some years before it was found experimentally, though it can also be explained within the framework of classical physics (Cabannes (1928), Rocard, Placzek); its great importance rests rather on the facility with which it can be applied to the study of molecules, and on the colossal amount of material relating to it which has been accumulated so quickly. The effect was discovered simultaneously (1928) by Raman in India, and by Landsberg and Mandelstam in Russia. They found that scattered light contains, in addition to the frequency of the incident light, a series of other frequencies.

The classical explanation of this effect is as follows. If a light wave  $E = E_0 \cos 2\pi\nu t$  falls on a molecule, it produces in this molecule a dipole moment

$$p = \alpha E_0 \cos 2\pi\nu t,$$

where  $\alpha$  denotes the polarizability tensor. As we have already remarked, the direction of the induced polarization is not in general the same as that of the exciting field strength. In the present case what we have to consider is the effect produced on the induced polarization by the state of the molecule as regards rotation and vibration. For, when the molecule rotates, the ellipsoid of polarization turns along with it, and the induced dipole moment therefore vibrates in the same rhythm. Similarly, when the nuclei oscillate, the whole electronic system also does so, and this again causes an oscillation of the polarizability in the same rhythm; in fact, since the electronic motion round the nuclei is much more rapid than the oscillations of the nuclei themselves, we can use the averaged electronic distribution when we are considering the effect of the nuclear vibrations on the polarizability. From the Fourier series which represents the influence of the rotations and vibrations on the polarizability we pick out a single term  $\nu_s$ , and write  $\alpha$  in the form

$$\alpha = \alpha_0 + \alpha_1 \cos(2\pi\nu_s t + \delta);$$

here  $\delta$  denotes an undetermined phase which varies from molecule to



# VIII] BAND SPECTRA AND THE RAMAN

molecule. If we substitute this expression in the form moment, we find

$$\begin{aligned}\mathcal{P} &= \alpha_0 \mathbf{E}_0 \cos 2\pi \nu t + \alpha_1 \mathbf{E}_0 \cos 2\pi \nu t \cos(2\pi \nu_s t + \delta) \\ &= \alpha_0 \mathbf{E}_0 \cos 2\pi \nu t \\ &\quad + \tfrac{1}{2} \alpha_1 \mathbf{E}_0 \{ \cos [2\pi(\nu + \nu_s)t + \delta] + \cos [2\pi(\nu - \nu_s)t - \delta] \}.\end{aligned}$$

The vibration of the dipole moment may therefore be regarded as due to the superposition of three vibrations with the frequencies  $\nu$ ,  $\nu + \nu_s$  and  $\nu - \nu_s$ . Since the dipole moment is the cause of the scattered light, the latter contains the frequencies  $\nu + \nu_s$  and  $\nu - \nu_s$  in addition to the frequency  $\nu$  of the incident light; moreover, since the phase  $\delta$  is arbitrary, the three vibrations are incoherent. Precisely the same holds good for all the rotational and vibrational frequencies of the molecule.

The spectrum of the scattered light is therefore a sort of band in the neighbourhood of the incident line, from which the rotational and vibrational frequencies can be read off, exactly as in the case of emission and absorption bands. The advantage of the method, however, is that the whole band is situated at a part of the *visible spectrum* which can be chosen at will, its position depending only on the choice of the irradiating frequency. Thus, for example, we can determine the constants of the ground state of the electronic system by observations, in the visible spectral region, of the separations of the lines of the bands, while the corresponding emission bands lie far in the infrared. It should be noted that the displacement of the bands into the visible, by observation of the rotational-vibrational bands in the case when a simultaneous electronic jump is involved, does not give the molecular constants of the unexcited "natural" state, but those of some excited state, which in general does not interest the chemist.

With regard, however, to the intensity relations in the Raman lines, special features are found, which do not admit of explanation on classical lines; for example, the component of the scattered light with the frequency  $\nu - \nu_s$  is much stronger than the one with the frequency  $\nu + \nu_s$ . This phenomenon can be understood at once in the light of the following simple quantum consideration. If an incident light quantum hits the molecule, it may as a first possibility be scattered without loss of energy. It may, on the other hand, excite the molecule, and so give up to it the vibrational energy  $h\nu_s$ ; in the scattered light the light quantum then appears with the energy  $h(\nu - \nu_s)$ . In rare cases, again, it may happen that a light quantum

hits a molecule which is already excited; it may then take energy away from the molecule, the latter falling into the lower state, while the light quantum flies on with energy  $h(\nu + \nu_s)$ . Hence on the long wave side the Raman lines are strong, but only a few weak lines lie on the short wave side.

The Raman effect has come to be of the greatest importance for molecular research, since it is comparatively simple to observe, and makes it possible to reach exact results regarding molecular structure in many cases, from the mere consideration that certain lines occur or do not occur. The following example serves to illustrate these remarks.

The molecule  $N_2O$ , assuming rectilinear arrangement, might have either of the forms NON or NNO. The question is how we can distinguish between these alternatives. Now the first form is obviously symmetrical, the second not. But, among the vibrations taking part in the Raman effect, those belonging to a symmetrical molecular structure differ characteristically (in the number and polarization of the lines) from those belonging to one which is unsymmetrical. Thus we can deduce without ambiguity from experiment that the non-symmetric form NNO is the correct one. Similar considerations can also be adduced with respect to complicated molecules (for instance, the symmetric structures methane,  $CH_4$ , and carbon tetrachloride,  $CCl_4$ ), but we cannot go into this here.

#### 4. Chemical Binding. Classification of Types of Binding.

Hitherto we have considered the electronic system of the molecule as a whole, and confined our attention in the main to properties and effects in relation to which the molecule is regarded as a given structure. We must now proceed to deal with the question which to the chemist is the specially interesting one, viz. how a molecular binding can come about at all, or how our molecular model is produced from the individual separate atoms.

We distinguish several different kinds of chemical binding, but between the principal sorts all possible intermediate stages occur. In our classification we adhere to the distinction drawn by Franck, according to which the decisive criterion is whether a molecule in dissociating tends more readily to split up into ions or into neutral atoms. It is true that the type of dissociation sometimes depends also on the excitation level of the molecule; still Franck's criterion gives in all cases a suitable point of view for the classification of binding. We therefore distinguish the following cases:

I. *The molecule splits up into ions* in dissociation by electrolysis

or in consequence of absorption of radiation; this kind of binding is called *ionic binding*.\* The extreme case of this type of binding is that in which the atoms are charged even in the molecule, so that they are present as ions; the binding is then explained by their Coulomb attraction. In this case we speak of *polar* (or *heteropolar*) binding. The typical case of this kind of binding is that of NaCl. About a hundred years ago it was conjectured by Berzelius that all chemical forces are really of electrostatic nature. The hypothesis was dropped, however, on account of the difficulty of explaining in this way the binding of atoms of the same kind (e.g.  $H_2$ ,  $N_2$ , . . .), which cannot possibly be of polar nature. It was only after it had become possible, by observation of electrolytic behaviour, of dipole moments and the like, to demarcate polar molecules in some measure from others, that the hypothesis of Berzelius was revived; it holds good for a limited class of bindings only, but within this range, as has been shown by Lewis and Kossel (1916), is capable of yielding important results.

II. *The molecule breaks up into atoms*, as, for instance, in the case of thermal excitation, or in consequence of absorption of radiation; we speak in this case of *atomic binding*. Here there are several sub-cases to be distinguished:

1. Atomic bindings with saturation of valency: *valency bindings* (also called *homopolar bindings*); to these belong in the first place the diatomic gases, such as  $H_2$ ,  $N_2$ ,  $O_2$ , as well as most organic compounds, for instance,  $CH_4$ .

2. Loose bindings without saturation of valency, due to the van der Waals forces: *cohesion bindings*.

3. Bindings which are effective, for example, in the case of the lattice formation of metals, and which we comprise under the name *metallic bindings*.

4. A series of bindings which do not admit of inclusion in the above scheme, such as the *benzene binding*, the *diamond binding*, and other similar binding types.

Between these various groups there is of course a series of intermediate stages, which will not, however, be considered here, the next sections being confined to a discussion of the most important of the types of binding mentioned above.

\* The nomenclature is not quite fixed, but that given here seems to be gradually becoming established. As a matter of history, it may be noted that the distinction between heteropolar and homopolar binding was first drawn by Abegg; the separation into ionic and atomic binding, with the criterion of the dissociation products, was introduced by Franck.

### 5. Theory of Heteropolar Ionic Binding.

How this binding comes about, we have already indicated above; we shall now consider the question a little more fully (Kossel, Born and Landé, 1918). The atoms which come next the inert gases are always striving to become converted into the inert gas configuration, by taking up or giving up electrons. In the alkali atoms the valency electron moves outside a closed shell and is comparatively loosely bound—the alkalis have a very small *ionization energy*  $I$  (see the numerical values of the ionization potential in Table IV, pp. 163, 164, 165). Conversely, in the halogens there is one electron too few to make up a closed shell, which, as we know, represents a very stable electronic configuration; they are therefore very ready to pick up an electron in order to complete the shell. We call this the *electronic affinity*  $E$ ; it is given by the energy which is set free when the electron settles in its place. Strictly connected with these concepts is that of *electrovalency*; positive electrovalency being the number of electrons loosely bound and therefore outside a closed shell in the atom, and negative electrovalency the number of electrons required to complete the inert gas configuration.

The genesis of an ionic binding may be pictured as a process consisting of two steps. In the first, the two reacting atoms become charged in opposite senses; thus, for example,  $\text{Na} + \text{Cl} = \text{Na}^+ + \text{Cl}^-$ . The second step is the attraction between the two ions in accordance with Coulomb's law; the energy of the attraction is  $-e^2/r$ . By itself, this would lead to the absolute coalition of the two ions; however, at small distances repulsive forces come into play, which can be accounted for by quantum mechanics. To represent them, a law of the form  $b/r^n$  has been tried, with good success; quantum mechanics gives approximately an exponential law  $be^{-r/\rho}$ , which has been found to answer even a little better. The position of equilibrium is then given by the value of  $r$  for which the sum of this part of the energy and the Coulomb energy has a minimum value.

In the gas molecule, in consequence of the one-sided action of the electric forces, the electron clouds of the two ions are of course very much deformed, so that the analysis of the ionic binding becomes extremely difficult. The circumstances are simpler in crystals, especially in the highly symmetrical ones of the rock salt (cubic face-centred, fig. 9) and similar types. In these the deformation disappears on account of the symmetry; in a rock salt lattice the same force is exerted on four sides on a chlorine ion by the neighbouring sodium ions, and a similar result holds for the more distant ions. Thus

we can calculate the *lattice energy*, that is, the energy  $U$  which must be supplied in order to break up the lattice completely. It is given by a sum of the form

$$U = \Sigma \left( \pm \frac{e^2}{r} + \frac{b}{r^n} \right),$$

the summation being taken over all the lattice points; here we adopt the same hypothesis as before with regard to the repulsive forces. The evaluation of this sum presents considerable difficulties, since the first or Coulomb part is only slowly convergent; the acting forces certainly diminish as the distance increases, but on the other hand the number of ions, at the same distance from a given ion on which they act, increases as the square of the distance. Advantageous methods of evaluating such lattice sums have been given by Madelung and Ewald (1918); the expression found for  $U$  has the form

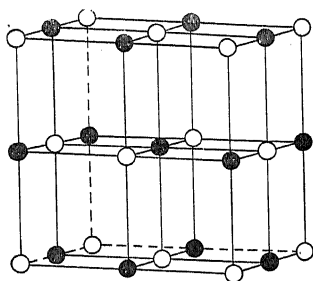


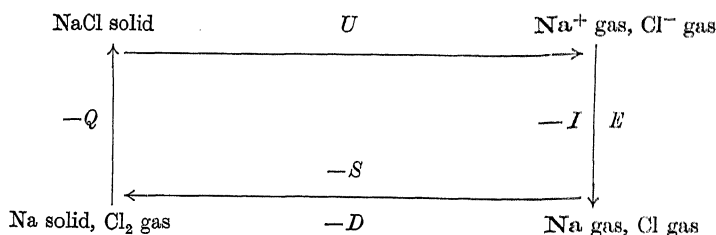
Fig. 9.—Face-centred cubic lattice

$$U = -\frac{e^2\alpha}{\delta} + \frac{\beta}{\delta^n},$$

where  $\delta$  is the lattice constant, i.e. the distance of a sodium ion from the nearest chlorine ion;  $\alpha$ , which is called Madelung's constant, depends on the lattice type, and for NaCl has the value 1.75. To determine the other constants,  $\beta$  and  $n$ , which are unknown initially, two equations are available. There is first the equation of equilibrium  $dU/d\delta = 0$ , expressing that the lattice energy is a minimum; secondly,  $d^2U/d\delta^2$  is the force needed to compress the crystal by a certain amount, and so can be determined experimentally. For the alkaline halides values of  $n$  between 6 and 10 are obtained in this way. Better results are found by using the exponential law  $be^{-r/\rho}$  for the repulsive force, as given by wave mechanics; the constant  $\rho$  is found to have approximately the same value for all the alkaline halides, viz.  $\rho \sim 0.35 \text{ \AA.}$  (Born and Mayer, 1932).

The results predicted by the theory can be tested directly. Thus, an experimental determination of lattice energies by J. Mayer, by means of thermal dissociation, showed good agreement for several salts. It is also possible to test the theory indirectly by calculating

the *electronic affinity* of a halogen derived from different salts (say of Cl from the compounds LiCl, NaCl, &c.), with the help of the following cyclic process (Born, 1919; Haber):



We start with solid rock salt (crystalline). By adding the lattice energy  $U$  (energy supplied is to be estimated as positive), we break up the crystal into ions. By further addition of the energy  $E$  corresponding to the electronic affinity of Cl, we can remove the excess electron from the  $\text{Cl}^-$ , whereupon the  $\text{Na}^+$  ion picks this electron up, at the same time giving up the ionization energy  $I$ , and so forming neutral Na gas. This gives up the heat of sublimation  $S$  of metallic Na, and is converted into solid sodium, while the atomic Cl gas gives up the dissociation energy  $D$  of the halogen and is converted into molecular  $\text{Cl}_2$  gas. By the action of this gas on the Na metal, crystalline rock salt is formed, the heat of formation  $Q$  being given up; and thus the cycle is closed. The energy of course must balance, so that we must have

$$U - I + E - S - D - Q = 0.$$

Here all the quantities are known except the electronic affinity  $E$ , viz.  $I$ ,  $D$ ,  $S$ ,  $Q$  from thermal and electrical measurements, and  $U$  from the lattice theory; and we can therefore calculate  $E$  from the

TABLE XI.—Electronic affinity  $E$  of Cl from the alkaline chlorides (in kcal per mole)

LiCl	NaCl	KCl	RbCl	CsCl	Mean
85.7	86.5	87.1	85.7	87.3	86.5

TABLE XII.—Electronic affinities of the halogens (in kcal per mole)

F	Cl	Br	I
95.3	86.5	81.5	74.2

equation. In the result, the same value must be found from all salts of the same halogen, and this, to a good approximation, is found to be the case. From the chlorine salts, for example, the values shown in the accompanying Table XI were found for the electronic affinity of Cl. The average values found for the different halogens are given in Table XII.

Recently, Mayer has succeeded in confirming these results by direct measurement, by means of molecular rays. It may be mentioned further, that for ions in solutions of salts there exists an absorption spectrum, which is referable to this process of liberation of the electron from the halogen ion; by taking into consideration the action of the surrounding medium (water), the electronic affinity is verified to a rough approximation (Franck and Scheibe).

## 6. Theory of Valency Binding.

We proceed now to consider the so-called valency binding. Here (according to Lewis, 1923) experience shows that such binding is frequently associated with the existence of an *electron pair*, which is such that each of the pair is shared by, or belongs to both of, the two atoms which are combined with each other. The hydrogen molecule is to be regarded as the simplest case of this kind, being built up from two equal nuclei and two electrons.

A method of dealing with the problem of the hydrogen molecule by quantum mechanics was pointed out by London and Heitler (1927). The following practical electrical example is exactly analogous to the case we are to consider, and may help to make the treatment of the general case easier to follow. If two like electrical oscillating circuits are brought close to each other, the original frequency of each being  $\nu_0$ , the coupling throws them to some extent out of tune,  $\nu_0$  being split up into two different frequencies, one higher than  $\nu_0$  and the other lower, so that beats are produced by the combination. The conditions in the hydrogen molecule are analogous to those mentioned; the two oscillating circuits are represented by the electrons revolving round the nuclei (in separate atoms). The coupling puts them out of tune a little, and we thus get a frequency somewhat higher than  $\nu_0$  and one somewhat lower. But, as we know, to every frequency  $\nu_0$  there corresponds an energy  $h\nu_0$ ; hence the undisturbed energy  $2E_0$  of the two separate hydrogen atoms gives rise to a somewhat lower and a somewhat higher energy of the coupled system

$$E_1 = 2E_0 - W_1(R), \quad E_2 = 2E_0 + W_2(R),$$

where  $W_1(R)$  and  $W_2(R)$  denote practically the coupling energies, which depend on the distance  $R$  between the two atoms, i.e. the nuclear distance (see Appendix XXXI, p. 340). The state represented by  $E_1$  possesses less energy than the dissociated state (separate hydrogen atoms), and therefore corresponds to a binding;  $E_2$  denotes repulsion of the two atoms (see fig. 7, p. 243). The energy  $W_1(R)$  corresponds

to the frequency with which the deformation of the wave functions by the interaction oscillates from one atom to the other. Since the square of the amplitude of a wave function represents the probability of finding a particle at a given place, this can be interpreted as an "exchange" of the two electrons. Therefore  $W_1(R)$  is called "exchange energy".

The phenomenon of *valency saturation* is explained by *Pauli's exclusion principle*. Just as in the case of atoms, so also in that of molecules the atoms, in consequence of this principle, become arranged in shells round the two nuclei. In the hydrogen molecule, the two electrons are in the ground state in the innermost "shell"; they must therefore, like the two electrons of a *K* shell, have opposite spins. Thus, in the attraction-state  $E_1$  the spins of the two electrons are antiparallel; on the contrary, it may be shown that in the state  $E_2$  the spins would be parallel. Here again "closed shells" represent specially stable states, and a chemical binding will easily come about if, in the molecule, saturation of the spins of the outer electrons occurs, that is to say, if a closed shell is formed. Generally, only those electrons in an atom, which do not already form an antiparallel pair with another electron, co-operate to a first approximation (the atomic distance  $R$  being great) in the binding. The number of spins not compensated is equal to the number of free valencies of the molecule (atom). The *true (homopolar) valency* is therefore equal to the *number of electrons with spins not saturated*. As examples we take the first two rows of the periodic table. In the first row we have H and He. Univalent hydrogen we have already discussed. Helium possesses two spins already saturated in a closed shell, and is therefore chemically neutral. The first element in the second row, Li, is, like hydrogen, univalent; it has one radiating electron outside a closed shell. With Be another electron comes in, whose spin becomes saturated in presence of the first *L* electron; Be should accordingly, like He, have valency 0. Actually it has valency 2, the explanation possibly being that it is only after excitation that the directions of spin in the two electrons come to have the same sense. Then follow the three elements B, C, N with 1, 2, 3 electrons respectively in the outermost shell ( $2p$  electrons). If a fourth  $p$  electron enters, so forming O, it will necessarily become saturated by means of one of those already present; for, when a definite spin direction is prescribed, there are only three possibilities for the setting of the orbital angular momentum. The last corner of the electrons must therefore, by the exclusion principle, take its place with the opposite spin direction, and so compensates one of the three



free spins of the nitrogen configuration of the electrons; oxygen is therefore divalent. When we pass on to fluorine, a new electron is added, which again saturates one of the two spins still present, making fluorine univalent. Neon with its closed shell is of course chemically neutral—its chemical valency is 0. This general view is well confirmed on the experimental side. We have already spoken of the difficulty with regard to Be; it must also be remembered that the beryllium compounds are mostly heteropolar. A more serious objection to the theory is that it ascribes two valencies to carbon, whereas this element is well known to be quadrivalent. This difficulty has also received a simple explanation. The carbon atom possesses, close above the ground state, an excitation level, in which it is quadrivalent; the ordinary chemical carbon bonds are due therefore to the first excited state of the carbon atom.

For the sake of brevity we shall not enter here into the mathematical treatment of valency binding (see Appendix XXXI, p. 340); we may mention, however, as special *successes of the theory* that it makes it possible to deal quantitatively with nuclear distances, heats of dissociation, and nuclear vibrational frequencies for molecular hydrogen. In other diatomic molecules also, the theory has proved capable of giving qualitative and quantitative results, which are in good agreement with experiment.

The interpretation of polyatomic compounds has been found difficult. Here the semi-empirical theories of Herzberg, Hund, Slater, Pauling and others pave the way to a satisfactory theory. A typical example of the phenomena to be explained is the fact, which has long been well known to physicists, that the three atoms of the water molecule do not lie in a straight line, as one might at first expect, but form an isosceles triangle (fig. 10). This phenomenon appears to have been successfully explained, at least qualitatively.

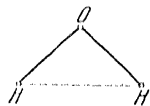


Fig. 10.—Model of water molecule

## 7. Theory of van der Waals Forces and other Kinds of Binding.

Of this subject we shall give only a very brief summary. The van der Waals cohesion forces depend on the mutual deformation of the atoms, and that in two different ways. First, the action of the field which emanates from the permanent dipole or quadrupole of the molecule, on the dipole induced in the other molecule by that field, leads on the average to an attraction; a result which, even before the quantum theory, had been proved by Debye and Keesom (1921) by classical considerations. From this, however, it would follow that

spherically symmetrical atoms (as of the inert gases) or molecules should show no cohesion, which is contrary to the experimental fact that all gases can be condensed. A way out of this difficulty was shown by London (1930), who proved that the deformability has a second effect, which is characteristic of the quantum theory. According to this theory there exists in fact a "zero-point motion", i.e. even in the lowest energy state an atom or molecule has to be put in correspondence with a system of moving charges (electrons), so that it carries a dipole oscillating with electronic frequency. If then two such systems come near each other, the zero-point motions of the dipoles always act in such a way that on the average the result is attraction; and the calculation gives an interaction energy which is inversely proportional to the sixth power of the nuclear distance, or  $W \sim 1/R^6$ . In many atoms or molecules without free valencies (spins), even before condensation a sort of pairing occurs, which is attributed to the van der Waals forces of attraction; it has to do with comparatively loose molecules with small dissociation energy. Examples of this sort of binding are given by metallic molecules, such as  $\text{Hg}_2$  and others. At sufficiently low temperatures, in all gases liquefaction occurs, and finally solidification. By means of his theory of cohesion, London (Appendix XXXII, p. 343) was able, to a fair approximation, to calculate the heats of sublimation of molecular lattices from atomic properties (ionization energy, proper frequency).

The existence of metals cannot be explained by the types of binding dealt with so far; we must regard them as being so formed that the positive atomic residues are held together by the free metallic electrons. A few cases, lithium, for example, have been successfully worked out (Slater, Wigner). Special difficulties are presented by a type of lattice (diamond lattice) in which the binding apparently depends on the saturation of the (four) valencies; among crystals of this type there occur both metals (Ge, Sn) and non-conductors, the latter including the extremely hard diamond itself (besides SiC, AgI, &c.). The relations involved have been elucidated by Hund, from the point of view of quantum mechanics.

Other examples, which in each case require the application of special theoretical considerations, are the ring formations of organic chemistry, the benzene ring, for example (Hückel); according to chemists' ideas the 6 carbon atoms are arranged in the form of a hexagon, with the hydrogen atoms outside (fig. 11). In order to secure that the carbon appear as quadrivalent in the structural diagram, the chemist has to introduce double bonds as shown. This unsatis-

factory lack of symmetry in the binding form is avoided by quantum mechanics, according to which the binding is effected by means of electrons, which revolve in the hexagon of carbon atoms.

### 8. Conclusion.

We have now reached the final stage of our task. In the preceding chapters we have in the main confined our account to the *positive results of research*, and have not dealt in detail with problems which

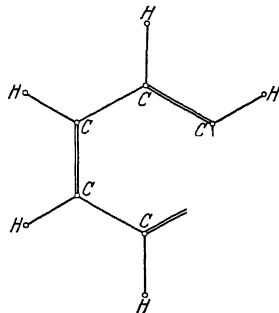


Fig. 11.—Benzene ring. Chemists assume double bonds, in order to allow for the quadrivalence of carbon.

are not yet completely cleared up. The impression might thus be given that physics had attained to a certain degree of finality. That, however, is by no means the case. There still remain many perplexing questions, all strictly connected with one another, and all in the last resort hanging upon the question of the structure of the nucleus. The questions are: which are the real elementary particles, electrons, positrons, protons or neutrons? Are photons, possibly also neutrinos, to be added to the list, or is the part they play a different one, perhaps by reason of the fact that their rest mass is nil? Why are the other particles of two kinds, one kind with small mass, the other with large? What are the forces which retain these particles in the nuclei? What are the laws of the stability and instability of nuclei?

Many physicists expect the answers to these questions to come from a blend of the theory of relativity and the quantum theory. The methods in use at present relate practically to the one body problem only, and fail when applied to processes involving several particles in rapid relative motion. The essential point here is the dualistic idea, according to which particles with their de Broglie waves, and electromagnetic fields, are regarded as two entirely different things. On the contrary, the author is firmly convinced that a satisfactory theory must be unitary in character, i.e. it must assume only one carrier of physical processes, the electromagnetic field, of which the particles are in some sense singularities. This conviction is based above all on the fact that particles, electrons and positrons, can be annihilated, becoming converted into electromagnetic waves; and that conversely they can be born in pairs from such waves. Dirac's theory of the positron (considered as a "hole" in a multitude of electrons with

negative energy) is a brilliant description of these processes, but not a satisfying explanation. All nuclear processes are transformations and separations of particles, which are so closely allied to the elementary process of amalgamation just referred to, that one is strongly inclined to refer them to the same root cause. In the author's opinion, the suggested generalization of Maxwell's electrodynamics, with the finiteness of the energy of point charges preserved, defines the direction in which research should proceed.

We look to the physics of the future for the solution of the enigma of inorganic matter. But the principal results of modern physics reach far beyond the domain in which they are won. As Niels Bohr first pointed out, the new views with regard to causality and determinism, which have arisen as a result of the quantum theory, are also of great significance for the biological sciences and for psychology. If even in inanimate nature the physicist comes up against absolute limits, at which strict causal connexion ceases and must be replaced by statistics, we shall be prepared, in the realm of living things, and emphatically so in the processes connected with consciousness and will, to meet insurmountable barriers, where mechanistic explanation, the goal of the older natural philosophy, becomes entirely meaningless. But this has the effect of completely changing the philosophical import of research. Physicists of to-day have learned that not every question about the motion of an electron or a light quantum can be answered, but only those questions which are compatible with Heisenberg's principle of uncertainty. There is a hint here for the biologist and the psychologist, that they should search for natural limits to causal explanation in their domains also, and mark out those limits with the same exactness as the quantum theory is capable of doing by means of Planck's constant  $h$ .

This is a programme of modesty, but at the same time one of confident hope. For what lies within the limits is knowable, and will become known; it is the world of experience, wide, rich enough in changing hues and patterns to allure us to explore it in all directions. What lies beyond, the dry tracts of metaphysics, we willingly leave to speculative philosophy.

## APPENDIX

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### I. Evaluation of Some Integrals Connected with the Kinetic Theory of Gases (p. 15).

Integrals of the type

$$I_\nu = \int_0^\infty v^\nu e^{-\lambda v^2} dv,$$

where  $\lambda = m\beta/2 = m/2kT$ , frequently occur in the kinetic theory of gases. The general form of the integral may at once be obtained from the particular cases  $I_0$  and  $I_1$  by differentiation with respect to  $\lambda$ . Thus, for example,

$$\begin{aligned} I_2 &= -\frac{dI_0}{d\lambda}, & I_3 &= -\frac{dI_1}{d\lambda}, \\ I_4 &= +\frac{d^2I_0}{d\lambda^2}, & I_5 &= +\frac{d^2I_1}{d\lambda^2}. \end{aligned}$$

$I_1$ , the second of the two fundamental integrals, can be evaluated by elementary methods:

$$I_1 = \int_0^\infty v e^{-\lambda v^2} dv = \frac{1}{2\lambda}.$$

$I_0$  is Gauss's well-known probability integral

$$I_0 = \int_0^\infty e^{-\lambda v^2} dv = \frac{1}{2} \sqrt{\frac{\pi}{\lambda}}.$$

The values of the succeeding integrals obtained, as stated above, by differentiation are as follows:

$$\begin{aligned} I_2 &= \int_0^\infty v^2 e^{-\lambda v^2} dv = \frac{1}{4} \sqrt{\frac{\pi}{\lambda^3}}, \\ I_3 &= \int_0^\infty v^3 e^{-\lambda v^2} dv = \frac{1}{2\lambda^2}, \end{aligned}$$

$$I_4 = \int_0^\infty v^4 e^{-\lambda v^2} dv = \frac{3}{8} \sqrt{\frac{\pi}{\lambda^5}},$$

and so on.

*Application to the integrals occurring in the kinetic theory of gases.*—

For  $n$ , the total number of molecules, we have

$$n = 4\pi A \int_0^\infty v^2 e^{-\frac{1}{2}\beta m v^2} dv = 4\pi A I_2 = A \sqrt{\frac{\pi^3}{\lambda^3}}.$$

Similarly, the total energy  $E$  is equal to

$$4\pi A \int_0^\infty \frac{1}{2} m v^4 e^{-\frac{1}{2}\beta m v^2} dv = 2\pi A m I_4 = \frac{3}{4} m A \sqrt{\frac{\pi^3}{\lambda^5}}.$$

The expressions for  $A$  and  $\beta$  given in the text (p. 15) are obtained by combining these two formulæ.

Integrals of this type also occur in the calculation of mean values. As according to Maxwell the number of molecules having a velocity between  $v$  and  $v + dv$  is

$$n_v dv = 4\pi n \left( \frac{m}{2\pi kT} \right)^{\frac{3}{2}} v^2 e^{-\lambda v^2} dv,$$

the “mean velocity”  $\bar{v}$  is given by

$$\bar{v} = \frac{\int_0^\infty n_v v dv}{\int_0^\infty n_v dv} = \frac{I_3}{I_2} = \frac{2}{\sqrt{(\pi\lambda)}}.$$

Again, the “mean square velocity”  $\sqrt{(\bar{v}^2)}$ , which is frequently used in the kinetic theory of gases, is given by

$$\bar{v}^2 = \frac{\int_0^\infty n_v v^2 dv}{\int_0^\infty n_v dv} = \frac{I_4}{I_2} = \frac{3}{2\lambda}.$$

$\sqrt{(\bar{v}^2)}$  is a little larger than  $\bar{v}$ , their ratio being  $\sqrt{(3\pi/8)} = 1.085$ . With these quantities we may compare the “most probable velocity”  $v_p$ , which corresponds to the maximum of the Maxwell distribution curve. It is given by the equation  $\frac{dn_v}{dv} = 0$ , or

$$\frac{d}{dv} (e^{-\lambda v^2} v^2) \equiv 2v e^{-\lambda v^2} (1 - \lambda v^2) = 0;$$

hence

$$v_p = \frac{1}{\sqrt{\lambda}}.$$

In the Maxwell distribution curve given in the text (fig. 5, p. 15), the three values  $v_p$ ,  $\bar{v}$ , and  $\sqrt{\bar{v}^2}$  are shown in their order and approximately proper positions.

Multiplying the expression for  $v_p$  above and below by the square root of Avogadro's number  $L$  and remembering that  $Lk = R$  and  $Lm = \mu$ , we have

$$\begin{aligned} v_p &= \sqrt{\frac{2kT}{m}} = \sqrt{\frac{2RT}{\mu}}; \\ \bar{v} &= \frac{2v_p}{\sqrt{\pi}} = \sqrt{\frac{8RT}{\pi\mu}}; \\ \sqrt{\bar{v}^2} &= v_p \sqrt{\frac{3}{2}} = \sqrt{\frac{3RT}{\mu}}. \end{aligned}$$

## II. Heat Conduction, Viscosity, and Diffusion (p. 19).

In this appendix we shall give a combined account of three phenomena by means of which the mean free path in a gas can be determined experimentally. These phenomena are *heat conduction*, *viscosity*, and *diffusion*. In all these phenomena there is a variation in some physical property of the molecules of the gas from point to point, which, however, tends to disappear as a result of the movements of the molecules.

*Heat conduction* occurs when external conditions of any kind give rise to a temperature gradient in a gas, i.e. when the molecules of gas at different points of space have different mean kinetic energies. Heat transference takes place owing to molecules from the warmer regions moving into cooler regions and giving up their surplus energy there, while slower-moving molecules move into the warmer regions and diminish the kinetic energy of the faster molecules there.

The circumstances are similar in the case of *viscosity*. This manifests itself in the form of a resistance acting on the faster-moving parts of the gas under consideration. According to the kinetic theory of gases, this resistance is due to molecules from slower-moving regions moving into faster-moving regions; as they then have a smaller mean velocity of flow than their surroundings, they will on the average be accelerated as a result of collision with the surrounding molecules, while the latter will be retarded, i.e. will be subject to a

resistance. We shall not go into the experimental methods for determining viscosity, but merely mention the result (due to Maxwell), which perhaps at first sight is surprising: the viscosity, like the thermal conductivity, is within wide limits independent of the pressure of the gas.

The third phenomenon which we shall consider here is *diffusion*. If in a mixture of two gases the concentration of one gas varies from point to point, i.e. if the ratio of the concentrations of the two gases differs at different points of the region occupied by the gaseous mixture (it is assumed that the pressure is everywhere the same, i.e. that the total number of gas molecules is the same), then it is clear that the molecules of one gas will gradually move from the regions where the concentration of this gas is greatest into regions where it is less; and similarly for the molecules of the other gas.

These three phenomena can be discussed mathematically in a very simple way provided we confine ourselves to a qualitative survey. In order to treat them together, we assume that the property  $A$  varies from point to point of space, e.g. has a non-vanishing derivative  $dA/dz$  in the direction of the  $z$ -axis. Then for heat conduction  $A$  is to be taken as the mean kinetic energy of a single gas molecule, for viscosity as the mean velocity of translation of the molecule in the direction of flow, and for diffusion as the number of molecules of a particular gas in a cubic centimetre of gaseous mixture. Now this variation of the property  $A$  from point to point gives rise to a transference of the property; a definite number of molecules cross unit area normal to the  $z$ -axis per second in either direction, and this number is given, at least approximately, by the product  $n\bar{v}$ , where  $n$  is the number of molecules per cubic centimetre and  $\bar{v}$  their mean velocity. But the molecules crossing the surface from one side possess the property  $A$  to a greater or less extent than those crossing the surface from the other side, so that there is a transference of the property across the surface. The quantity  $M(A)$  crossing per second can easily be estimated if we note that any molecule has a mean free path  $l$  between two successive collisions with other molecules, i.e. in the interval of time between successive collisions it describes on the average a path whose length is of the order  $l$ ; for the purposes of this approximation we are not concerned with the exact numerical factor. For  $M(A)$  we readily obtain the expression

$$M(A) \sim n\bar{v}\{A(z_0 - l) - A(z_0 + l)\},$$

where  $z_0$  is the co-ordinate of the element of area considered. Expanding, we have (apart from numerical factors)



$$M(A) \sim -n\bar{v}l \left( \frac{dA}{dz} \right)_{z_0}.$$

The amount of the property  $A$  transferred is therefore proportional to the "gradient of  $A$ ", and also to the number of molecules per cubic centimetre, their mean velocity, and their mean free path. This equation is known as the *transport equation*.

We note that  $M(A)$  is independent of the pressure, provided that  $A$  itself denotes a property of the gas molecules which is independent of the pressure (Maxwell). For the pressure of a gas is given by  $p = nkT$ , i.e. at constant temperature depends only on the number of molecules per cubic centimetre. It is true that  $n$  appears as a factor in the transport equation; but this factor is compensated by the occurrence in the formula of the mean free path, which is inversely proportional to  $n$  and to the cross-section of the molecule. This independence of the pressure accordingly results from the fact that though more molecules take part in the transference of  $A$  at the higher pressures, they do not on the average travel so far.

We shall now make particular application of the transport equation to the three phenomena mentioned above. We begin with *heat conduction*. Here  $A$  stands for the kinetic energy of a molecule, i.e.  $E_{\text{kin}} = \text{const.} + c_v mT$ , where  $c_v m$  is the specific heat of the gas at constant volume, for a single molecule ( $c_v$  is the specific heat per mole). The quantity of heat  $Q$  which crosses unit area per second is then given by

$$Q \sim -n\bar{v}l c_v m \frac{dT}{dz}.$$

We see that it is proportional to the temperature gradient; the factor of proportionality  $\kappa = n\bar{v}l c_v m$  is called the *thermal conductivity*.

In the case of the *viscosity*, as we saw above,  $A$  stands for  $mu$ , the mean linear momentum of a molecule resulting from the flow of the gas. Then the momentum transferred per second (per unit area of the surface of contact between faster-moving portions of gas and slower-moving), i.e. the resistance  $R$ , is given by

$$R \sim -n\bar{v}l m \frac{du}{dz};$$

the quantity  $\eta = n\bar{v}l m$  is called the *coefficient of viscosity*. We see that the quotient  $\kappa/\eta c_v$  is a constant of the order of unity. Theoretically this constant must be the same for all molecules of the same

structure, i.e. the quotient must have a constant value for all monatomic gases, another constant value for all diatomic gases, and so on. Our qualitative discussion of course does not enable us to obtain the exact value of the constant.

As we have already emphasized, the discussion given here is of course only a rough sketch. Improvements and refinements in the theory have been made by Boltzmann, Maxwell, and others, by considering the mechanism of collision and the distribution of velocities in greater detail; these improvements, however, yield no new principle, but merely lead to greater accuracy in the numerical factors. Here, however, we shall not go into the matter further. It remains for us to point out that the above theory is not valid unless the mean free path is small compared with the dimensions of the vessel containing the gas. If this is not the case (at atmospheric pressure  $l$  is of the order of  $10^{-6}$  cm., but is equal to about 10 cm. at the pressure in an X-ray tube ( $10^{-4}$  mm. of mercury)), the laws which hold are quite different. The molecules then fly practically straight from one wall of the containing vessel to the other without colliding with other molecules. Thus, for example, if there is a difference of temperature between two opposite parts of the wall of the container, they carry heat energy directly from one wall to the other. The quantity of heat transferred is then proportional to the number of molecules;  $n$  no longer disappears from the transport equation, as the mean free path no longer enters into it. The laws of heat conduction, viscosity, &c., at low pressures have been especially studied by Knudsen. They are of great practical importance, for instance, in connexion with the working of air pumps (such as Gaede's rotary molecular pump and the diffusion pump).

We shall now briefly consider the problem of *diffusion*. We imagine a mixture of two gases in dynamical equilibrium, i.e. the pressure, and hence  $n$ , the total number of molecules per cubic centimetre, are to be the same throughout. Here the property  $A$  is  $n_1/n$ , the concentration of one kind of molecule, or  $n_2/n$ , the concentration of the other kind of molecule. Then the transport equation gives the number ( $Z_1$ ) of molecules of the first kind, or the number ( $Z_2$ ) of molecules of the second kind, that diffuse through unit area in unit time:

$$Z_1 \sim -n\bar{v}l \frac{d(n_1/n)}{dz} = -\bar{v}l \frac{dn_1}{dz}, \quad Z_2 \sim -\bar{v}l \frac{dn_2}{dz}.$$

If the phenomenon is a steady one,  $n_1 + n_2$  must be equal to a constant  $n$ , i.e.  $dn_1/dz = -dn_2/dz$ ; then the total flow  $Z_1 + Z_2$  is zero.

The two kinds of molecule have the same diffusion constant  $\delta = \bar{v}l$ , which in virtue of the factor  $l$  is inversely proportional to  $n$ , the total number of molecules.

### III. Van der Waals' Equation of State (p. 20).

As compared with the equation of state for perfect gases, van der Waals' equation of state for actual gases, given in the text (p. 20), contains two correction terms, a volume correction and a pressure correction. Here we shall seek to show, at least qualitatively, how these terms arise.

(1) The fact that in the equation of state, as is stated in the text, exactly four times the total volume of the molecules themselves must be subtracted from the total volume of the gas, may be explained as follows. In § 6 (p. 9) we investigated the probability that  $n$  molecules should be distributed in a given way among the cells  $\omega_1, \omega_2, \dots$ . We found it to be the product of the number of ways in which a definite distribution, prescribed by the numbers  $n_1, n_2, \dots$  of the molecules occupying the individual cells, can be realized and the a priori probability of the occurrence of this distribution. It is with this a priori probability that we are concerned here. If we inquire into the probability that  $n$  molecules will be found in a definite portion of volume  $v$ , we begin by assuming, as was done in § 6 (p. 9), that the probability is proportional to  $v^n$ . This is assuredly the case so long as we can neglect the finite magnitude of the gas molecules, as, e.g., in rarefied gases. It is not so, however, for high pressures, where the gas molecules are so tightly packed together that their own volume is actually comparable with that of the space available for them. Here we obtain the desired result in the following way. Let  $v_m$  be the volume of a molecule (i.e. for a spherical molecule  $v_m = \frac{4}{3}\pi(\frac{1}{2}\sigma)^3$ , where  $\sigma$  denotes the diameter of the molecule). Now the centres of two molecules cannot approach within a distance equal to the diameter of a molecule; hence each molecule has an effective volume of magnitude  $\frac{4}{3}\pi\sigma^3 = 8v_m$ , independent of the particular shape which the molecule happens to have.

The probability that a molecule will be found in a definite portion of volume  $v$  is of course proportional to  $v$ , as above. If we introduce a second molecule into this region, the space available for it is only  $v - 8v_m$ ; the space left for a third molecule is  $v - 2 \times 8v_m$ , and so on. The probability of finding  $n$  molecules in  $v$ , then, is proportional, not to  $v^n$ , but to the product

$$v(v - 8v_m)(v - 2 \times 8v_m) \dots (v - (n - 1)8v_m);$$

$v$  is accordingly replaced by the  $n$ -th root of this product. We can easily calculate this; as  $v_m$  is very small and  $8v_m n$  still small in comparison with  $v$ , we can replace the product approximately by

$$v^n - 8v^{n-1}v_m(1 + 2 + \dots + (n-1)) \sim v^n \left(1 - \frac{8v_m}{v} \frac{n^2}{2}\right).$$

Taking the  $n$ -th root of this, we see that if the finite volume of the molecules is to be taken into account  $v$  must be replaced by

$$v \left(1 - n \frac{4v_m}{v}\right) = v - b,$$

where  $b$  stands for  $4nv_m$ , four times the actual volume of the molecules contained in  $v$ .

(2) The term  $a/v^2$  added to the pressure may be explained as follows. If there are forces of cohesion acting between the molecules, one element of volume acts on another of equal size with a force which is proportional to  $n^2$ , where  $n$  is the number of molecules per cubic centimetre. The pressure exerted by the gas on an external body is correspondingly less than it would be if there were no forces of cohesion. Hence in the equation of state  $p$  must be replaced by  $p + An^2$ . Now let  $nV = N$  be the total number of molecules in the volume of gas  $V$ ; this number of course does not vary if the volume is changed, so that the additional term  $An^2$  may be written in the form  $A \left(\frac{N}{V}\right)^2 = \frac{a}{V^2}$ . The constant  $a$  accordingly is closely related to the energy of cohesion of the gas, and we may mention that the magnitude of the latent heat of evaporation of liquids also depends on it. Van der Waals' equation is still approximately true even for densities which correspond to the liquid state, and enables the phenomenon of liquefaction to be investigated thermodynamically.

#### IV. The Mean Square Deviation (p. 21).

All phenomena involving deviation from a mean value depend on the formula

$$\overline{\Delta n^2} = \bar{n}.$$

Here  $n$  may denote, e.g., the number of particles in a definite fixed portion of volume of a gas. If we stick to this example, we know that this number is not always the same but varies with the time. There is, however, a time-average ( $\bar{n}$ ) of  $n$ , about which the number of particles

varies. If we could observe the actual number  $n$  at any instant (as if in a snapshot), we should obtain varying values  $n_1, n_2, \dots$ , deviating from the mean value  $\bar{n}$  by the amounts  $\Delta n_1 = n_1 - \bar{n}, \Delta n_2 = n_2 - \bar{n}, \dots$ . The sum of these deviations for a large number of observations, divided by their number, must of course vanish; a result differing from zero, however, is obtained if we average the squares of these deviations over a large number of observations. We thus obtain the mean square deviation  $\overline{\Delta n^2}$ , which according to the assertion made above is equal to the mean value  $\bar{n}$ .

To prove this, we start from the fundamental formulæ of the kinetic theory of gases, according to which the probability that of  $N$  molecules in a total volume  $V$  the fraction  $n$  will be found in the portion of volume  $v$  is given by the formula

$$W_n = \frac{N!}{n!(N-n)!} \left(\frac{v}{V}\right)^n \left(\frac{V-v}{V}\right)^{N-n}$$

(cf. p. 12: instead of the distribution among the cells  $\omega_1, \omega_2, \dots$ , we here have the distribution between the two regions  $v$  and  $V-v$  with "occupation numbers"  $n_1 = n$  and  $n_2 = N-n$ ). The sum of all the probabilities of the various distributions is  $\sum_{n=0}^N W_n = 1$  (by the binomial theorem). The mean number ( $\bar{n}$ ) of molecules in  $v$  is obtained by working out the sum

$$\bar{n} = \sum_{n=0}^N n W_n.$$

If we put  $v/V = x$ , this sum becomes

$$\begin{aligned} \bar{n} &= \sum_{n=0}^N \frac{nN!}{n!(N-n)!} x^n (1-x)^{N-n} \\ &= Nx \sum_{n=1}^N \frac{(N-1)!}{(n-1)![(N-1)-(n-1)]!} x^{n-1} (1-x)^{(N-1)-(n-1)}. \end{aligned}$$

According to the binomial theorem the sum here is again equal to 1, so that for  $\bar{n}$ , the mean number of particles, we obtain the expression

$$\bar{n} = Nx = \frac{Nv}{V},$$

as might be expected. In order to calculate  $\bar{n}^2$ , we note that the mean value  $n(n-1)$  may be found in exactly the same way as  $\bar{n}$ . In fact,

$$\begin{aligned}\overline{n(n-1)} &= \sum_{n=0}^N n(n-1)W_n \\ &= N(N-1)x^2 \sum_{n=2}^N \frac{(N-2)!}{(n-2)![(N-2)-(n-2)]!} x^{n-2}(1-x)^{(N-2)-(n-2)},\end{aligned}$$

whence

$$\overline{n(n-1)} = N(N-1)x^2.$$

Hence we at once have

$$\overline{n^2} = \overline{n(n-1)} + \bar{n} = N(N-1)x^2 + Nx,$$

so that

$$\overline{\Delta n^2} = \overline{(n - \bar{n})^2} = \overline{n^2} - \bar{n}^2 = Nx - Nx^2 = N \frac{v}{V} \left(1 - \frac{v}{V}\right).$$

If we confine ourselves to small values of  $x$ , i.e. to relatively small volumes  $v$ , we at once obtain the deviation formula

$$\overline{\Delta n^2} = \bar{n}$$

given above. Many other mean values can be calculated in exactly the same way as  $\bar{n}$  and  $\overline{n^2}$ .

These formulæ are applied as follows. Spontaneous deviations of the density of the molecules from the mean value are associated with a change in almost all the physical properties of the gas. For example, variations of density in a gas are associated with variations in the refractive index. By observing the variations of measurable properties we can determine the mean square deviation. For example, the variation of the refractive index gives rise to a scattering of the transmitted light, which is proportional to the mean square deviation (to this, according to Lord Rayleigh, is due the blue colour of the sky). In the Brownian movement, too, fine particles in suspension execute certain motions as a result of variations in the density of the surrounding medium, and these motions are a measure of the mean square deviation of the density of the surrounding medium. If we know the way in which the phenomenon in question is related to the density, the mean value of the number of particles ( $\bar{n}$ ) can be determined by measuring the deviations and calculating the mean square deviation.

An important field of application of the theory of deviations is the subject of radioactivity. Here it is a question of counting the number of particles emitted by a radioactive preparation per second (e.g. by means of a Geiger counting apparatus (p. 32)). If we are dealing with a long-lived substance, for which the average number ( $\bar{n}$ ) of

particles emitted per second may be regarded as constant, the actual number ( $n$ ) of particles emitted in individual periods of one second will differ from  $\bar{n}$ . The fact that the deviations *are* in accordance with the law stated above,  $\overline{\Delta n^2} = \bar{n}$ , forms a convincing proof of the statistical nature of the disintegration process. Corresponding formulæ can also be deduced for short-lived radioactive substances.

## V. Theory of Relativity (p. 27).

In classical mechanics it is proved that an observer who experiments only within a closed system cannot determine whether this system is at rest or is in uniform motion. In fact, the Newtonian equations of motion  $m d^2x/dt^2 = F$  (where  $m$  is the mass,  $F$  the force,  $x$  the co-ordinate of a particle, and  $t$  the time) remain unchanged if we pass to a moving co-ordinate system by the transformation  $x' = x - vt$ , provided the force depends only on the position of the particle relative to the co-ordinate system (since  $x_1' - x_2' = x_1 - x_2$ ). This principle of relativity in mechanics has to be modified when electromagnetic pro-

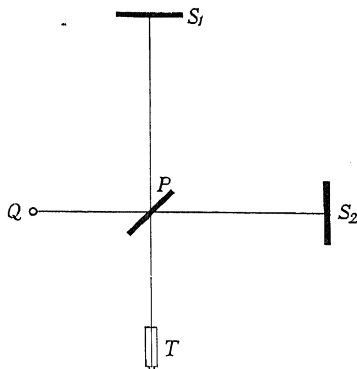


Fig. 1.—Diagram of Michelson's interferometer.  $Q$ , source of light;  $P$ , semi-transparent, silvered glass plate;  $S_1$ ,  $S_2$ , mirrors;  $T$ , telescope.

cesses, light waves for example, are taken into account. Since these waves move in vacuo, it was the custom to assume a carrier, the æther. The earth moves through it, so that an observer on the earth must perceive an "æther wind", which retards or accelerates the light waves according to their direction. The experiment was carried out by means of Michelson's interferometer (fig. 1). A light ray from the source  $Q$  was divided into two parts by partial reflection at the thinly silvered glass plate  $P$ ; the two partial rays were reflected at the mirrors  $S_1$ ,  $S_2$  and united again at the plate  $P$ . Interference fringes are seen in the telescope  $T$ . If the apparatus is then turned round so that first  $PS_1$ , then  $PS_2$ , fall in the direction of the æther wind, the interference fringes ought to be displaced. The result of the experiment was negative: the æther wind is not really there.

To explain this fact Einstein developed his theory of relativity.

The leading idea is that the customary combination of space and time—kinematics—must be abandoned. There is no absolute time, but just as every moving system has its “proper” co-ordinates  $x, y, z$ , so it has also a proper time  $t$ , which has to be transformed as well as the co-ordinates when we pass to a new system. The equations defining this so-called Lorentz transformation for two systems moving in the  $x$ -direction with the relative velocity  $v$  are

$$x' = \frac{x - vt}{\sqrt{(1 - v^2/c^2)}}, \quad y' = y, \quad z' = z, \quad t' = \frac{t - vx/c^2}{\sqrt{(1 - v^2/c^2)}}$$

where  $c$  is the velocity of light. It is easily verified that these equations give

$$x'^2 - c^2 t'^2 = x^2 - c^2 t^2.$$

This formula suggests that  $x, y, z, ict$  ( $i = \sqrt{-1}$ ) may be interpreted (Minkowski) as co-ordinates in a four-dimensional space, in which  $x^2 + y^2 + z^2 + (ict)^2$  represents the square of the distance from the origin; a Lorentz transformation then represents a rotation round the origin in this space. Minkowski's idea has developed into a geometrical view of the fundamental laws of physics, culminating in the inclusion of gravitation in Einstein's so-called general theory of relativity.

Physically, the equation last written expresses the fact that  $x = ct$  implies  $x' = ct'$ ; or that the velocity of light is independent of the motion of the observer. The negative result of Michelson's experiment is thus explained.

Further, we see that if the distance of two points is  $x_1 - x_2$  at the same time  $t$  in one system  $(x, t)$ , then in a second system  $(x', t')$  their distance is

$$x_1' - x_2' = (x_1 - x_2)/\sqrt{(1 - v^2/c^2)}.$$

Thus

$$x_1 - x_2 = (x_1' - x_2')\sqrt{(1 - v^2/c^2)}.$$

Hence lengths in the second system appear from the first system to be shortened (the Fitzgerald-Lorentz contraction). On the other hand, if  $t_1, t_2$  are the times of two events at the same place  $x$  in the first system, then

$$t_1' - t_2' = (t_1 - t_2)/\sqrt{(1 - v^2/c^2)},$$

so that the time between the two events seems longer from the second system (cf. p. 271).

It must be assumed that there is no velocity greater than the velocity of light—otherwise the theory would become meaningless.



It follows that the basic ideas of mechanics must be so altered that the motion of a body can never be accelerated up to the velocity of light.

We can arrive at this result from the consideration that the velocity defined by the components  $dx/dt$ ,  $dy/dt$ ,  $dz/dt$  (or the momentum obtained from this velocity on multiplication by the mass) cannot, in view of Lorentz's transformation, be regarded as a vector, since the differential  $dt$  in the denominator is also transformed. We obtain a serviceable "covariant" definition if we replace  $dt$  by  $dt_0$ , where  $dt_0$  is the element of the "proper" time of the particle, i.e. the time measured in the system of reference in which the particle is at rest. The relation between  $dt$  and  $dt_0$  is found by taking the derivative of  $t'$ ,

$$\frac{dt'}{dt} = \frac{1 - \frac{v}{c^2} \frac{dx}{dt}}{\sqrt{(1 - v^2/c^2)}},$$

and putting  $dx/dt = v$ ;  $dt'$  then becomes  $dt_0$ , and we have

$$dt_0 = \sqrt{(1 - v^2/c^2)} dt.$$

The component momentum is now defined to be

$$m_0 \frac{dx}{dt_0} = \frac{m_0}{\sqrt{(1 - v^2/c^2)}} \frac{dx}{dt} = m \frac{dx}{dt},$$

where  $m_0$  is a constant, the rest mass. The mass  $m$  is therefore given by the formula of the text,  $m = m_0/\sqrt{(1 - v^2/c^2)}$ . The formula has been confirmed, not only by the experiments with cathode rays already mentioned (p. 27), but also by certain details in the behaviour of spectral lines, particularly those of hydrogen. In fact, since such lines are emitted by rapidly moving electrons, they reflect the mechanical properties of these electrons (p. 107). The energy is

$$E = mc^2 = \frac{m_0 c^2}{\sqrt{(1 - v^2/c^2)}},$$

and the momentum is

$$G = mv = \frac{m_0 v}{\sqrt{(1 - v^2/c^2)}}.$$

From these two expressions we find

$$m_0 = \frac{1}{c^2} \sqrt{(E^2 - G^2 c^2)},$$

the relation which is used in Chap. III, § 4, p. 59, in the discussion on

the existence of the neutrino. The values of  $E$  and  $G$  therefore determine the rest mass. If the latter vanishes, as it does for a light quantum (photon, § 2, p. 69),  $E$  and  $G$  are connected by the relation

$$cG = E.$$

## VI. Electron Theory (p. 44).

Maxwell's equations, for an isotropic substance, are:

$$\frac{1}{c} \frac{\partial \mathbf{D}}{\partial t} - \text{curl } \mathbf{H} = -\frac{4\pi}{c} \mathbf{i}, \quad \text{div } \mathbf{D} = 4\pi\rho, \quad \mathbf{D} = \epsilon \mathbf{E};$$

$$\frac{1}{c} \frac{\partial \mathbf{B}}{\partial t} + \text{curl } \mathbf{E} = 0, \quad \text{div } \mathbf{B} = 0, \quad \mathbf{B} = \mu \mathbf{H}.$$

We assume the reader to be acquainted with the vector notation, and the meaning of the vectors  $\mathbf{E}$ ,  $\mathbf{H}$ ,  $\mathbf{D}$ ,  $\mathbf{B}$ ,  $\mathbf{i}$ , as also with the energy theorem

$$-\frac{dW}{dt} = \int S_n d\sigma + A,$$

where  $W = (1/8\pi) \iiint (\epsilon \mathbf{E}^2 + \mu \mathbf{H}^2) dx dy dz$  is the electromagnetic energy;  $A = \iiint \mathbf{i} \mathbf{E} dx dy dz$  is the rate of doing work by the electric force on the currents; and  $\mathbf{S} = (c/4\pi)[\mathbf{E}\mathbf{H}]$  is the Poynting vector, of which use is made in Appendix VIII, p. 275.

Lorentz's theory of the electron considers fields in vacuo only, and therefore puts  $\epsilon = \mu = 1$ ; the current is assumed to be a convection current,  $\mathbf{i} = \rho \mathbf{v}$ , the charge being rigidly bound to the electron. To connect the theory with mechanics, it is further necessary to assume that the field exerts the mechanical force

$$\mathbf{F} = \int \int \int \rho \left\{ \mathbf{E} + \frac{1}{c} [\mathbf{v} \mathbf{H}] \right\} dx dy dz$$

on the electron.

In the author's new field theory Maxwell's equations are retained, but  $\epsilon$  and  $\mu$  are given by the equations\*

$$\mu = \frac{1}{\epsilon} = \sqrt{1 + \frac{1}{b^2} (\mathbf{B}^2 - \mathbf{E}^2)}$$

(in the text, for brevity,  $\mathbf{H}$  is written in place of  $\mathbf{B}$ ); the

\* We consider here the theory in its original form; it has been modified by adding a term  $\frac{1}{b^4} (\mathbf{E}\mathbf{B})^2$  under the square root, which has no influence on the statical solution.

electrons are taken to be point charges ( $\rho = i = 0$ ). The energy is

$$W = \frac{1}{4\pi} \iiint \{ \epsilon E^2 + b^2(\mu - 1) \} dx dy dz,$$

the Poynting vector

$$\mathbf{S} = \frac{c}{4\pi} [\mathbf{D}\mathbf{B}] = \frac{c}{4\pi} [\mathbf{E}\mathbf{H}].$$

In the statical case we have  $\mathbf{H} = \mathbf{B} = 0$ , so that

$$\mathbf{D} = \epsilon \mathbf{E} = \frac{\mathbf{E}}{\sqrt{\left(1 - \frac{1}{b^2} E^2\right)}}.$$

For a point singularity, it follows from  $\text{div } \mathbf{D} = 0$  that

$$D = \frac{e}{r^2},$$

and hence that

$$E = \frac{e}{\sqrt{(r^4 + a^4)}}, \quad a^2 = \frac{e}{b}.$$

The field  $\mathbf{E}$  therefore remains finite even at the centre of the electron, where it has the value

$$E_0 = \frac{e}{a^2} = b.$$

Since the field equations reduce to  $\text{curl } \mathbf{E} = 0$ , a potential  $\phi$  exists, which is found from  $E = -\partial\phi/\partial r$  to be

$$\phi = \frac{e}{a} \int_{r/a}^{\infty} \frac{dx}{\sqrt{(1+x^4)}}.$$

For  $r \gg a$ ,  $\phi$  is easily seen to reduce to the Coulomb potential  $e/r$ ; but for  $r = 0$  it is finite,

$$\phi(0) = 1.8541 \frac{e}{a}.$$

If we substitute the expression for  $\mathbf{E}$  in the energy formula

$$W = \frac{1}{4\pi} \iiint \left\{ \frac{E^2}{\sqrt{(1 - E^2/b^2)}} + b^2(\sqrt{1 - E^2/b^2} - 1) \right\} dx dy dz$$

we find

$$W = \frac{2}{3} e \phi(0) = 1.2361 \frac{e^2}{a}.$$

Einstein's theorem  $W = mc^2$  is here rigorously true, but the proof would take us too far. Equating  $mc^2$  to the value of  $W$  given above, we get  $a = 1.2361 e^2/mc^2$ , which, on introducing the empirical values for the charge and the mass of the electron, comes out as  $a = 2.28 \times 10^{-13}$  cm. The absolute field is  $b = e/a^2 = 9.18 \times 10^{15}$  e.s.u. It can be shown that the equations of motion given by Lorentz in the classical theory,

$$\frac{d}{dt}(mv) = e(E + \frac{1}{c}[\mathbf{v}H]),$$

hold here rigorously for a constant external electromagnetic field, and approximately for a field which does not vary appreciably over the "diameter of the electron"  $2a$ . For quickly vibrating fields there are deviations, namely a diminution of the force. It may be that this result can be applied to the explanation of the extremely high penetrating power of cosmic radiation (§ 7, p. 41), which is much less absorbed than the usual theories predict.

The adaptation of these ideas to the principles of quantum theory and the introduction of the spin has been carried out to some extent, but still meets with great difficulties.

## VII. The Theorem of the Inertia of Energy (p. 47).

Here we give the mathematical proof corresponding to the ideal experiment given in the text. Classical electrodynamics, in agreement with experiment, yields the result that the momentum transferred by light of energy  $E$  to a surface absorbing it is  $E/c$  (cf. Appendix XXVII, p. 329); hence the momentum transferred by recoil to the box during the emission is of the same magnitude. If the box (fig. 1, p. 46) has the total mass  $M$ , it acquires a recoil velocity  $v$  to the left, given according to the theorem of the conservation of momentum by  $Mv = E/c$ . The box continues to move during the time required by the light to traverse the distance ( $l$ ) between the transmitter ( $I$ ) and the receiver ( $II$ ). If we neglect terms of higher order, this time  $t$  is equal to  $l/c$ . During this time the box moves a distance  $x = vt = El/Mc^2$  to the left. In order not to contradict the fundamental principle of the centre of gravity we must assume, as was explained in the text, that the transference of energy from  $I$  to  $II$  is accompanied by a simultaneous transference of mass in the same direction. If we denote this meanwhile unknown mass by  $m$ , then the total moment due to displacement of mass after the phenomenon is given by  $Mx - ml$ ; by the fundamental principle of the centre of gravity this must be zero.

Substituting for  $x$  the value found above, we obtain Einstein's formula for  $m$ :

$$m = M \frac{x}{l} = \frac{E}{c^2}.$$

### VIII. Calculation of the Coefficient of Scattering for Radiation of Short Wave-length (p. 48).

In order to calculate the coefficient of scattering we start from the familiar formulæ for the radiation from a dipole antenna. According to Hertz, the field of a dipole of moment  $\mathcal{P}$  is given by

$$|E| = |H| = \frac{|\ddot{\mathcal{P}}|}{c^2 r} \sin \theta.$$

The electric vector is at right angles to the magnetic vector and both are at right angles to the direction in which the radiation is propagated;  $\theta$  is the angle between the direction of propagation of the radiation and the direction in which the dipole is oscillating, and  $r$  the distance from the dipole. A necessary assumption for the validity of this formula is that  $r$  must be very much greater than the wave-length of the radiation emitted, or, in other words, the above formula is valid only for the so-called wave zone.

To apply the formula to our case of a vibrating particle, we note that this represents an oscillating electric dipole of moment  $\mathcal{P} = es$ , where  $s$  is the displacement of the electron at any instant from its position of equilibrium. Then the energy radiated per second in a specified direction is given by the Poynting vector, the length of which is

$$S = \frac{c}{4\pi} |E| |H| = \frac{e^2 |\ddot{s}|^2}{4\pi c^3 r^2} \sin^2 \theta.$$

The total energy radiated in unit time is obtained by integrating this over the whole surface of a sphere:

$$I = \int S d\sigma = \int_0^\pi \int_0^{2\pi} S r^2 \sin \theta d\theta d\phi.$$

As

$$\int \sin^2 \theta d\sigma = 2\pi r^2 \int_0^\pi \sin^3 \theta d\theta = \frac{8\pi r^2}{3},$$

the total radiation is given by

$$I = \frac{2}{3} \frac{e^2}{c^3} |\ddot{s}|^2.$$

This formula still contains the displacement of the oscillating electron, or, more accurately, its second derivative with respect to the time. This is determined by the equation of motion of the electron under the influence of the incident radiation:

$$m\ddot{\mathbf{S}} = e\mathbf{E}_0.$$

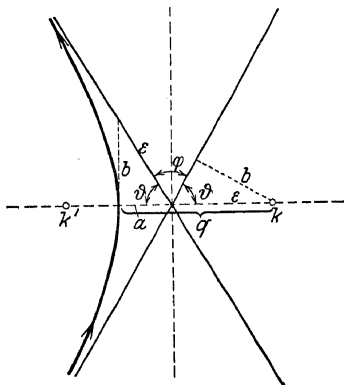
$\mathbf{E}_0$ , the electric vector of the primary wave, however, is connected with the intensity of the incident radiation by the equation

$$I_0 = \frac{c}{4\pi} |\mathbf{E}_0| |\mathbf{H}_0| = \frac{c}{4\pi} E_0^2.$$

If we substitute these expressions in the equation above, the formula for the energy received by the electron from the primary wave and transformed into scattered radiation becomes

$$I = \frac{2}{3c^3} \frac{e^4}{m^2} E_0^2 = \frac{8\pi}{3} \left( \frac{e^2}{mc^2} \right)^2 I_0.$$

As the mass of the scattering particle occurs squared in the denominator, we see that a proton or an atomic nucleus scatters some millions of times less than an electron.



### IX. Rutherford's Scattering Formula for $\alpha$ -rays (p. 53).

According to Rutherford, the nucleus of an atom (with charge  $Ze$ ) and an  $\alpha$ -particle (with charge  $E$  and

Fig. 2.—Hyperbolic path in the scattering of an  $\alpha$ -ray by a nucleus;  $a$  and  $b$  are the semi-axes,  $\epsilon$  the distance between the centre and the focus. The angle of deflection is  $\phi = \pi - 2\theta$ .

mass  $M$ ) repel one another with the Coulomb force  $ZeE/r^2$ . If we consider the heavy nucleus as at rest, the  $\alpha$ -particle describes an orbit (fig. 2) which is one branch of a hyperbola, one of whose foci coincides with the nucleus  $K$ . Let  $b$  be the distance of the nucleus from the asymptote of the hyperbola which would be described by the  $\alpha$ -particle if there were no repulsive force. Further, we denote the distance of the nucleus  $K$  from the vertex of the hyperbola by  $q$ ; then

$$q = \epsilon(1 + \cos \theta),$$

where  $\epsilon$  denotes the linear eccentricity  $OK$  (i.e. the distance between the centre  $O$  and the focus  $K$ ), and  $\theta$  the angle between the asymptote and the axis. From fig. 2 we readily see that

$$\epsilon = \frac{b}{\sin \theta}$$

and hence that

$$q = \frac{b(1 + \cos \theta)}{\sin \theta} = b \cot \frac{\theta}{2};$$

$b$  is obviously equal to the length of the minor semi-axis of the hyperbola.

In the first place we shall seek to find a relation between the "collision parameter"  $b$  and the angle of deviation  $\phi$ , which, as we see from the figure, is equal to  $\pi - 2\theta$ . We accordingly apply the laws of motion to the  $\alpha$ -particle. First we have the theorem of the conservation of energy: the sum of the kinetic energy and the potential energy is constant. At a very great distance from the nucleus the  $\alpha$ -particle has kinetic energy only; let its velocity there be  $v$ . If we equate this energy to the total energy at the instant when the electron is passing the vertex of the hyperbola, we have

$$\frac{1}{2}Mv^2 = \frac{1}{2}Mv_0^2 + \frac{ZeE}{q},$$

or, if we divide by  $\frac{1}{2}Mv^2$  and for short put  $k = ZeE/Mv^2$ ,

$$\frac{v_0^2}{v^2} = 1 - \frac{2k}{b} \frac{\sin \theta}{1 + \cos \theta}.$$

Further, we have the theorem of the conservation of angular momentum, in virtue of which

$$Mvb = Mv_0q,$$

or

$$\frac{v_0}{v} = \frac{b}{q} = \frac{\sin \theta}{1 + \cos \theta},$$

$$\left(\frac{v_0}{v}\right)^2 = \frac{\sin^2 \theta}{(1 + \cos \theta)^2} = \frac{1 - \cos \theta}{1 + \cos \theta}.$$

Substituting this value in the above equation and carrying out a few transformations, we have

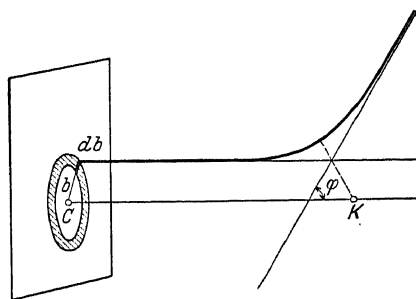
$$\frac{b}{k} = \frac{\sin \theta}{\cos \theta} = \tan \theta,$$

or, since  $\phi = \pi - 2\theta$ ,

$$b = k \cot \frac{\phi}{2}.$$

This gives the angle of deviation as a function of  $b$ , the distance of the rectilinear prolongation of the path of the  $\alpha$ -particle (the asymptote) from the nucleus.

It is now easy to calculate how many  $\alpha$ -particles in an incident parallel beam are deviated by a specified amount. We imagine a



plane  $E$  at right angles to the incident beam and at a great distance from  $K$ ;  $C$  is the foot of the perpendicular from  $K$  to  $E$  (fig. 3). It is obvious that all the  $\alpha$ -particles which pass through a ring of  $E$  formed by

Fig. 3.—Relative frequency, in a definite angular region, of scattering of  $\alpha$ -particles by a nucleus.

two circles with radii  $b$  and  $b + db$  will be subject to a deviation between  $\phi$  and  $\phi + d\phi$ . If one particle passes through one square centimetre of  $E$  in one second, the number passing through the ring in question will be

$$dn = 2\pi b db,$$

where

$$db = k d(\cot \phi/2) = -\frac{k d\phi}{2 \sin^2 \phi/2}.$$

Hence

$$|dn| = \pi k^2 \frac{\cos \phi/2}{\sin^3 \phi/2} |d\phi|.$$

This is the number of particles deviated through an angle between  $\phi$  and  $\phi + d\phi$ ; they are uniformly distributed over a zone of a unit sphere, the area of the zone being  $2\pi \sin \phi d\phi$ . Hence  $W(\phi)$ , the number of  $\alpha$ -particles deviated into a unit of area of the unit sphere, is  $dn/2\pi \sin \phi d\phi$ , so that the probability of deviation per unit solid angle is

$$W(\phi) = \frac{1}{4} k^2 \frac{1}{\sin^4 \phi/2} = \left( \frac{ZeE}{2Mv^2} \right)^2 \frac{1}{\sin^4 \phi/2}.$$

This is Rutherford's scattering formula. Each relationship (between



$Z, M, v, \phi$ ) contained in it can be tested experimentally by counting the scattered  $\alpha$ -particles; the  $v$ -relation, it is true, can only be investigated experimentally over a small region, as the range of velocity available with naturally-occurring  $\alpha$ -particles is small. In general the agreement between the formula and experimental results is extremely accurate; with light atoms, of course, we must take into account the recoil of the nucleus  $K$  after collision with the  $\alpha$ -particle, which is easily done. Marked discrepancies have been found only in the case of almost central collisions (deviations of almost  $180^\circ$ ) of light atoms (which have low nuclear charges, so that the incident  $\alpha$ -particle approaches the nucleus very closely); here, however, we shall not go into further details.

As the charge and mass of the  $\alpha$ -particles are known (they are  $\text{He}^{++}$  ions, for which  $M = 4M_H$ ,  $E = 2e$ ) and their velocities can be determined from deflection experiments, the formula can be used to find the nuclear charge  $Z$ . For this we need only know the number of scattering atoms per unit volume and count the number of  $\alpha$ -particles in the presence and absence of the scattering layer. For example, accurate experiments by Chadwick gave the following values for  $Z$ :

Platinum	Silver	Copper
77.4	46.3	29.3,

while according to the periodic table the numbers should be

78	47	29.
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The excellent agreement between the two sets of figures confirms the basic assumption that the nuclear charge and the atomic number are identical.

## X. The Compton Effect (p. 75).

Here we shall investigate the collision between a light quantum and an electron, on the assumptions of the special theory of relativity. This procedure is appropriate for our purpose, as it does not make the calculations any more complicated, while on the other hand the result obtained is then valid for the scattering even of very hard radiation.

The calculations are based on the theorems of the conservation of energy and of momentum. The energy of the light quantum before the collision is  $h\nu$ , its momentum  $h\nu/c$ ; the corresponding quantities after the collision we shall call  $h\nu'$  and  $h\nu'/c$ . For the sake of simplicity we shall regard the electron as at rest

previous to the collision; in this case its energy will be equal to the rest-energy  $m_0c^2$  corresponding by Einstein's formula to the rest-mass  $m_0$ , while its momentum will be zero. Let  $v$  be the velocity of the electron after the collision; then its mass  $m$  will be  $\frac{m_0}{\sqrt{1-v^2/c^2}}$ , its energy  $mc^2$  will be  $\frac{m_0c^2}{\sqrt{1-v^2/c^2}}$ , and its momentum  $mv$  will be  $\frac{m_0v}{\sqrt{1-v^2/c^2}}$ . We may also say that after collision the electron has the "kinetic energy"

$$(m - m_0)c^2 = m_0c^2 \left( \frac{1}{\sqrt{1-v^2/c^2}} - 1 \right)$$

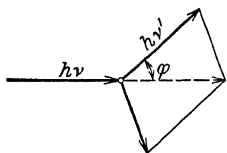


Fig. 4.—Momentum diagram for the Compton effect

(which, as we easily see by expanding in powers of  $v/c$ , agrees with the formula  $\frac{1}{2}m_0v^2$  of non-relativistic mechanics when the velocity is small), while before collision this "kinetic energy" is zero.

Then if  $\phi$  is the angle of deviation of the light quantum and  $\psi$  the angle of deviation of the electron, the theorems of the conservation of energy and of momentum take the following forms (see fig. 4):

$$\text{Conservation of energy: } h\nu + m_0c^2 = h\nu' + mc^2;$$

$$\text{Conservation of momentum: } \begin{cases} \frac{h\nu}{c} = \frac{h\nu'}{c} \cos \phi + mv \cos \psi, \\ 0 = \frac{h\nu'}{c} \sin \phi - mv \sin \psi. \end{cases}$$

Eliminating  $\psi$  from the last two equations, we obtain

$$m^2v^2c^2 = h^2\{(\nu - \nu' \cos \phi)^2 + (\nu' \sin \phi)^2\} = h^2(\nu^2 + \nu'^2 - 2\nu\nu' \cos \phi).$$

Again, from the energy equation we have

$$m^2c^4 = \{h(\nu - \nu') + m_0c^2\}^2 = h^2(\nu^2 + \nu'^2 - 2\nu\nu') + 2m_0c^2h(\nu - \nu') + m_0^2c^4.$$

$$\text{Since } m^2 = \frac{m_0^2}{1 - v^2/c^2}, \quad \text{i.e. } m^2 \left( 1 - \frac{v^2}{c^2} \right) = m_0^2,$$

by definition, subtracting one of these equations from the other gives

$$m_0^2c^4 = -2h^2\nu\nu'(1 - \cos \phi) + 2m_0c^2h(\nu - \nu') + m_0^2c^4$$

## APPENDIX XI

or 
$$(1 - \cos \phi) = \frac{m_0 c^2}{h} \frac{\nu - \nu'}{\nu \nu'} = \frac{m_0 c^2}{h} \left( \frac{1}{\nu'} - \frac{1}{\nu} \right).$$

If for convenience we denote the quantity

$$\frac{h}{m_0 c} = 0.0242 \text{ \AA.}$$

(which is in general referred to as the Compton wave-length) by  $\lambda_0$ , the above equation can also be written in the form

$$\Delta \lambda = \lambda' - \lambda = c \left( \frac{1}{\nu'} - \frac{1}{\nu} \right) = (1 - \cos \phi) \frac{h}{m_0 c} = 2\lambda_0 \sin^2 \frac{\phi}{2},$$

which was given in the text.

### XI. Phase Velocity and Group Velocity (p. 78).

In order to give a strict proof of the relationship  $U = \partial \nu / \partial \tau$  given in the text, we in the first instance consider the most general form of a group of waves; it must have the form of a Fourier integral

$$u(x, t) = \int a(\tau) e^{2\pi i(\nu t - \tau x)} d\tau,$$

where  $\nu = \nu(\tau)$  is to be regarded as a function of the wave-number  $\tau$ .

We now assume that the group is very narrow, so that in the integral there occur only those waves of finite amplitude whose wave-numbers differ from the mean wave-number  $\tau_0$  by a very small amount. If we put  $\tau = \tau_0 + \tau_1$ ,  $\nu(\tau) = \nu_0 + \nu_1(\tau_1)$ , and  $a(\tau_0 + \tau_1) = b(\tau_1)$ , the wave-group may be written in the form

$$u(x, t) = A(x, t) e^{2\pi i(\nu_0 t - \tau_0 x)},$$

where

$$A(x, t) = \int b(\tau_1) e^{2\pi i(\nu_1 t - \tau_1 x)} d\tau_1.$$

Hence the wave-group may be regarded as a single wave of frequency  $\nu_0$ , wave-number  $\tau_0$ , and amplitude  $A(x, t)$  varying from point to point and moment to moment. This assumption is justified, as according to our assumption  $A(x, t)$  is a function which varies only slowly compared with the exponential function  $e^{2\pi i(\nu_0 t - \tau_0 x)}$ ; to a first approximation it varies in the rhythm of a mean of the beat frequencies  $\nu_1$ , which are very small compared with  $\nu_0$ .

The velocity with which a definite value of the amplitude  $A(x, t)$ ,

e.g. its maximum, advances with the wave is called the group velocity. This is accordingly found from the relation

$$\frac{\partial A}{\partial x} \frac{dx}{dt} + \frac{\partial A}{\partial t} = 0$$

obtained by differentiating the equation  $A(x, t) = \text{const.}$  with respect to the time. If we call the group velocity  $U$  to distinguish it from the phase velocity, we have

$$U = \left( \frac{dx}{dt} \right)_{A=\text{const.}} = - \frac{\partial A / \partial t}{\partial A / \partial x}.$$

Now obviously

$$\frac{\partial A}{\partial t} = 2\pi i \int b(\tau_1) \nu_1 e^{2\pi i(\nu_1 t - \tau_1 x)} d\tau_1,$$

$$\frac{\partial A}{\partial x} = -2\pi i \int b(\tau_1) \tau_1 e^{2\pi i(\nu_1 t - \tau_1 x)} d\tau_1.$$

As we have assumed that the group is confined to a very narrow range of wave-length, we can expand  $\nu_1(\tau_1)$  in powers of  $\tau_1$ :

$$\nu_1(\tau_1) = \nu(\tau) - \nu_0 = \left( \frac{d\nu}{d\tau} \right)_0 \tau_1 + \dots$$

Hence

$$\frac{\partial A}{\partial t} = - \left( \frac{d\nu}{d\tau} \right)_0 \frac{\partial A}{\partial x},$$

and for the group velocity we accordingly have

$$U = \frac{d\nu}{d\tau},$$

while the phase velocity is given by

$$u = \frac{\nu}{\tau}.$$

## XII. Elementary Derivation of Heisenberg's Uncertainty Relation.

We consider a wave packet of finite width. For the sake of simplicity we represent its amplitude at any moment by a Gauss error function (such as actually occurs in the ground state of the harmonic oscillator, Appendix XXXII, p. 343):

$$f(x) = A e^{-x^2/a^2};$$

then  $\Delta x$  the width is given by

$$\Delta x = \sqrt{x^2} = \frac{1}{2}a.$$

The Fourier representation of  $f(x)$  is

$$f(x) = \int_{-\infty}^{\infty} \phi(\tau) e^{2\pi i \tau x} d\tau,$$

where

$$\phi(\tau) = \int_{-\infty}^{\infty} f(x) e^{-2\pi i \tau x} dx$$

is the amplitude of a partial harmonic wave with the wave number  $\tau$ . Introducing the expression of  $f(x)$  in  $\phi(\tau)$  we have

$$\begin{aligned} \phi(\tau) &= A \int_{-\infty}^{\infty} e^{-(x^2/a^2 + 2\pi i \tau x)} dx \\ &= A e^{-(\pi \tau a)^2} \int_{-\infty}^{\infty} e^{-(x/a + \pi i \tau a)^2} dx. \end{aligned}$$

This integral is transformed by the substitution  $x/a + \pi i \tau a = y$  into the Gauss integral

$$a \int_{-\infty}^{\infty} e^{-y^2} dy = a\sqrt{\pi}.$$

Therefore

$$\phi(\tau) = A a \sqrt{\pi} e^{-(\pi \tau a)^2} = A a \sqrt{\pi} e^{-\tau^2/b^2},$$

where

$$b = \frac{1}{\pi a}.$$

The distribution of the elementary waves composing the wave packet  $f(x)$  is again a Gauss function with the half width  $b = 1/\pi a$ . Putting

$\Delta x = \frac{1}{2}a$ ,  $\Delta \tau = \frac{1}{2}b$ , one has  $\Delta x \cdot \Delta \tau = \frac{1}{4\pi}$ , and introducing the momentum  $p = h\tau$ :

$$\Delta x \cdot \Delta p = \frac{h}{4\pi},$$

which is the exact expression of Heisenberg's uncertainty law for the special wave packet (see Appendix XXXII, p. 343). It is evident that with respect to order of magnitude the relation  $\Delta x \cdot \Delta p \sim h$  holds for any form of the wave packet. We shall prove the inequality with an exact numerical coefficient later (see Appendix XXII, p. 312).

### XIII. Hamiltonian Theory and Action Variables (p. 101).

Here we shall briefly consider the mechanics of multiply periodic motions and the corresponding quantum conditions. According to Hamilton, the motion of a system is described completely by stating the energy as a function of the co-ordinates  $q_k$  and the momenta  $p_k$ , the so-called Hamiltonian function  $H(q_1, q_2, \dots, p_1, p_2, \dots)$ . (For ordinary rectangular co-ordinates the momentum  $p_k$  becomes  $m_k \dot{q}_k$ ;

here, as in what follows, the dot denotes differentiation with respect to the time.) The equations of motion are then

$$\dot{q}_k = \frac{\partial H}{\partial p_k}, \quad \dot{p}_k = -\frac{\partial H}{\partial q_k}.$$

From these the theorem of the conservation of energy follows immediately; for if we form the total differential coefficient of  $H$  with respect to the time, and make use of the equations of motion, we have

$$\frac{dH}{dt} = \sum_k \left( \frac{\partial H}{\partial q_k} \dot{q}_k + \frac{\partial H}{\partial p_k} \dot{p}_k \right) = 0,$$

that is, the energy of the system,  $H(q_k, p_k) = E$ , is constant.

In general we can replace the pair of variables  $(p_k, q_k)$  by an arbitrary pair of canonically conjugate variables, where we define canonical variables by the fact that they satisfy equations of the type of the equations of motion given above. The investigation of such canonical transformations leads to a mathematical problem (the so-called Hamilton-Jacobi differential equation) which in many cases can be solved. Here we shall assume that a solution can be found. Then the problem of integrating the equations of motion may be stated in the following form: to find new canonical variables  $(J_k, w_k)$ , such that the energy depends on the quantities  $J_k$  only, not on the quantities  $w_k$ . Then by the equations of motion  $\dot{J}_k = -\frac{\partial H}{\partial w_k} = 0$ , so that  $J_k$  is constant throughout the motion. On the other hand,  $\dot{w}_k = \nu_k = \frac{\partial H}{\partial J_k}$  is also constant as time goes on, owing to the constancy of the quantities  $J_k$ , and  $w_k$  increases linearly with the time:  $w_k = \nu_k t$ . Thus the integration problem is solved in the new co-ordinates, and it only remains to make the reverse transformation.

A system is said to be singly or multiply periodic if the variables just defined can be found such that each rectangular co-ordinate is periodic in the quantities  $w_k$ , i.e. can be represented as a Fourier series in terms of these variables  $w_1, \dots, w_k$ ; so that,  $\tau_1, \tau_2, \dots$  being positive or negative integers,

$$x = \sum_{\tau_1, \tau_2, \dots} a_{\tau_1, \tau_2, \dots} e^{2\pi i(w_1 \tau_1 + w_2 \tau_2 + \dots)} = \sum_{\tau_1, \tau_2, \dots} a_{\tau_1, \tau_2, \dots} e^{2\pi i(\tau_1 \nu_1 + \tau_2 \nu_2 + \dots)t}.$$

The coefficients  $a_{\tau_1, \tau_2, \dots}$  are then functions of the quantities  $J_k$ . If the motion is periodic, the quantities  $J_k$  are said to be *action variables* and the quantities  $w_k$  *angle variables*. We have already had an example

of this kind in the rotator, where the azimuth  $\phi$  and the angular momentum  $p$  were taken as canonical co-ordinates. Then the Hamiltonian function is  $H = \frac{p^2}{2A}$ , whence there follow the equations of motion  $\phi = \frac{\partial H}{\partial p} = \frac{p}{A} = \omega$ ,  $\dot{p} = -\frac{\partial H}{\partial \phi} = 0$ . These have the solutions  $p = \text{const.}$  and  $\phi = \omega t$ . The rectangular co-ordinate is then given by

$$x = ae^{i\phi} = ae^{i\omega t}.$$

Ehrenfest has proved that the action variables  $J_k$  are adiabatic invariants, i.e. that the quantities  $J_k$  can be quantised. We now postulate the *quantum conditions*

$$J_k = hn_k \quad (n_k \text{ an integer}).$$

We can, however, only set up as many conditions of this kind as there are incommensurable frequencies in the motion. This may be seen as follows. By way of example we consider the case of two degrees of freedom; the index in the Fourier series then contains the sum  $\nu_1\tau_1 + \nu_2\tau_2$ , where  $\tau_1, \tau_2$  are integers. If, e.g.,  $\nu_1 = k\nu_2$ , where  $k$  is an integer,  $\nu_1\tau_1 + \nu_2\tau_2 = \nu_2(k\tau_1 + \tau_2)$ . Now  $k\tau_1 + \tau_2$  can take any integral value, since  $\tau_1$  and  $\tau_2$  are integers. In reality, therefore, we have only a single periodic motion and a single periodic Fourier series, and it is clear that in these circumstances only one quantum condition can exist. This case is referred to as *degeneracy*.

The canonical variables  $J_k$  and  $w_k$  can be conveniently determined in the case where the system is separable, i.e. once the equations of motion have been solved each  $p_k$  depends only on the corresponding  $q_k$ :

$$p_k = p_k(q_k).$$

Here it may be shown that the action variable  $J_k$  corresponding to the  $k$ th period is given by

$$J_k = \oint p_k dq_k,$$

where the integral is to be taken over the whole period. Hence in the case of separable co-ordinates the quantum condition may be stated at once:

$$\oint p_k dq_k = hn_k.$$

Without further examination these conditions do not reveal whether degeneracy is present or not. Hence before applying them it is necessary to investigate carefully what the actual number of incommensurable periods in the system is. Frequently, however, it is convenient

to take no notice of this (Sommerfeld) and simply to write down superfluous quantum conditions. Then the physically observable quantities like energy, momentum, &c., depend only on certain combinations of the quantum numbers  $n_1, n_2, \dots$ , so that they may be expressed by a smaller number of integers; thus contradictions cannot arise. This method has the following additional advantage: if a degenerate system is perturbed, e.g. by an external electric or magnetic field, the degeneracy is in general removed and there appears a new fundamental period, incommensurable with the former periods. If the unperturbed system has been discussed by means of suitable variables involving a superfluous period, and the perturbations are small, the quantum conditions can be taken over direct from the unperturbed case. We shall meet with an example of this in connexion with the Kepler problem.

#### XIV. Quantisation of the Elliptic Orbits in Bohr's Theory (pp. 92, 107).

The quantisation of the elliptic orbits in the Bohr atom is carried out here by means of the method, described in the previous section, of over-quantising the system by using more periods to describe the motion than it actually has.

We have first to write down the classical laws of motion for two particles of masses  $m$  and  $M$  and charges  $-e$  and  $Ze$  which attract one another according to Coulomb's law. (These correspond exactly to the laws of motion in astronomy, except that there the attracting force is the force of gravitation.) If  $(x_1, y_1, z_1)$  and  $(x_2, y_2, z_2)$  are the co-ordinates of the electron and the nucleus respectively,  $W$ , the energy of the motion, is given by

$$W = \frac{1}{2}m(\dot{x}_1^2 + \dot{y}_1^2 + \dot{z}_1^2) + \frac{1}{2}M(\dot{x}_2^2 + \dot{y}_2^2 + \dot{z}_2^2) - \frac{Ze^2}{r}.$$

If we regard the centre of gravity of the system as at rest, then  $mx_1 + Mx_2$  is zero; introducing the relative co-ordinates  $(x, y, z)$ , where  $x = x_2 - x_1$ , &c., as new variables, we have

$$x_1 = \frac{-M}{M+m}x, \quad x_2 = \frac{m}{M+m}x.$$

Similarly for the corresponding components of the velocity. The corresponding term of the kinetic energy is then

$$\frac{1}{2}m\dot{x}_1^2 + \frac{1}{2}M\dot{x}_2^2 = \frac{1}{2}\mu\dot{x}^2,$$

where  $\mu$  is the so-called effective mass  $mM/(m+M)$ ; cf. the discussion of the motion of the nucleus in § 1, p. 96. Hence the problem



may be treated as a Kepler problem for the relative motion about the centre of gravity with co-ordinates  $(x, y, z)$  and effective mass  $\mu$ .

For such motions about a fixed centre the conservation of angular momentum (theorem of areas) holds as well as the conservation of energy. The former implies in the first place that the motion is all in one plane. We take this plane as the  $xy$ -plane and transform to polar co-ordinates by means of the equations  $x = r \cos \phi$ ,  $y = r \sin \phi$ . The momentum theorem then gives

$$\mu r^2 \dot{\phi} = p_{\phi} = \text{const.}$$

In polar co-ordinates the energy is

$$W = \frac{1}{2} \mu (\dot{r}^2 + r^2 \dot{\phi}^2) - \frac{Ze^2}{r} = \text{const.}$$

Eliminating  $\dot{\phi}$  between these two equations and noting that

$$\dot{r} = \dot{\phi} \frac{dr}{d\phi} = \frac{p_{\phi}}{\mu r^2} \frac{dr}{d\phi},$$

we have

$$W = \frac{1}{2} \frac{p_{\phi}^2}{\mu} \left\{ \frac{1}{r^4} \left( \frac{dr}{d\phi} \right)^2 + \frac{1}{r^2} \right\} - \frac{Ze^2}{r}.$$

Here it is advantageous to introduce the new variable  $\rho = 1/r$ ; for  $\rho$  we then have the differential equation

$$W = \frac{1}{2} \frac{p_{\phi}^2}{\mu} \left\{ \left( \frac{d\rho}{d\phi} \right)^2 + \rho^2 \right\} - Ze^2 \rho.$$

This may readily be solved by differentiating once more with respect to  $\phi$ ; rejecting the factor  $d\rho/d\phi$ , we have the differential equation of the second order

$$\frac{d^2 \rho}{d\phi^2} + \rho - \frac{Ze^2 \mu}{p_{\phi}^2} = 0.$$

As is well known, its general solution is

$$\rho = \frac{Ze^2 \mu}{p_{\phi}^2} + C \cos(\phi - \phi_0),$$

containing the two constants  $C$  and  $\phi_0$ ; by suitable choice of  $\phi$  the latter can always be made to vanish. If for brevity we also put

$$q = \frac{p_{\phi}^2}{e^2 Z \mu}, \quad \epsilon = Cq = \frac{C p_{\phi}^2}{e^2 Z \mu},$$

we have

$$r = \frac{1}{\rho} = \frac{q}{1 + \epsilon \cos \phi}.$$

As we know, this equation represents a conic. In order that the motion may be periodic, i.e. in order that the conic may be entirely at a finite distance, the denominator in the above polar equation must never vanish; that is,  $|\epsilon|$  must be less than 1. Perihelion, i.e. the minimum value of  $r$ , corresponds to the azimuth  $\phi = 0$ ; then  $r_1 = \frac{q}{1 + \epsilon}$ . Aphelion corresponds to  $\phi = \pi$ :  $r_2 = \frac{q}{1 - \epsilon}$ . From these we obtain the expression

$$a = \frac{1}{2}(r_1 + r_2) = \frac{q}{1 - \epsilon^2}$$

for the major semi-axis  $a$ . For  $\phi = \pm \pi/2$ ,  $r$  is equal to the semi-latus-rectum  $q$  of the ellipse (see fig. 8, p. 107). Further, we readily see that  $\epsilon$  is the numerical eccentricity and that the minor semi-axis is given by  $b = \frac{q}{\sqrt{1 - \epsilon^2}}$ . The energy of the motion is obtained by substituting from the equation of the ellipse in the previous expression for  $W$ ; we have

$$\begin{aligned} W &= \frac{1}{2} \frac{p_\phi^2}{\mu} \left\{ \left( \frac{d\rho}{d\phi} \right)^2 + \rho^2 \right\} - Ze^2\rho \\ &= \frac{1}{2} \frac{p_\phi^2}{\mu} \left\{ C^2 + \left( \frac{Ze^2\mu}{p_\phi^2} \right)^2 \right\} - \frac{Z^2e^4\mu}{p_\phi^2} = -\frac{e^2Z}{2a}, \end{aligned}$$

which is the same as the result for a circular orbit (p. 92) if we replace the radius by the major semi-axis.

Now, although there is only one period, we prescribe two quantum conditions

$$\oint p_r dr = n'h, \quad \oint p_\phi d\phi = kh.$$

Owing to the constancy of  $p_\phi$ , the second condition at once gives

$$p_\phi = \frac{h}{2\pi} k.$$

The first integral, however, is not so easy to calculate; here we have

$$\begin{aligned} n'h &= \oint p_r dr = \oint \mu \dot{r} dr = \oint \mu \frac{dr}{d\phi} \dot{\phi} \frac{dr}{d\phi} d\phi \\ &= p_\phi \oint \frac{1}{r^2} \left( \frac{dr}{d\phi} \right)^2 d\phi. \end{aligned}$$

Accordingly, we have to evaluate the integral

$$I = \int_0^{2\pi} \frac{1}{r^2} \left( \frac{dr}{d\phi} \right)^2 d\phi = \int_0^{2\pi} \frac{1}{\rho^2} \left( \frac{d\rho}{d\phi} \right)^2 d\phi = \int_0^{2\pi} \frac{\epsilon^2 \sin^2 \phi}{(1 + \epsilon \cos \phi)^2} d\phi$$

for a complete revolution, i.e.  $\phi$  runs from 0 to  $2\pi$ . Integrating by parts, we have

$$I = \left[ \frac{\epsilon \sin \phi}{1 + \epsilon \cos \phi} \right]_0^{2\pi} - \int_0^{2\pi} \frac{\epsilon \cos \phi}{1 + \epsilon \cos \phi} d\phi = - \int_0^{2\pi} \frac{\epsilon \cos \phi}{1 + \epsilon \cos \phi} d\phi$$

Multiplying the second expression by  $+2$  and the first by  $-1$  adding, we obtain

$$\begin{aligned} I &= - \int_0^{2\pi} \frac{\epsilon^2 \sin^2 \phi + 2\epsilon \cos \phi (1 + \epsilon \cos \phi)}{(1 + \epsilon \cos \phi)^2} d\phi \\ &= - \int_0^{2\pi} \frac{1 + 2\epsilon \cos \phi + \epsilon^2 \cos^2 \phi}{(1 + \epsilon \cos \phi)^2} d\phi + \int_0^{2\pi} \frac{1 - \epsilon^2}{(1 + \epsilon \cos \phi)^2} d\phi. \end{aligned}$$

The first of these integrals can be evaluated at once and gives  $-2\pi$ ; the second, in virtue of the equation of the ellipse, may be written in the form

$$\int_0^{2\pi} \frac{1 - \epsilon^2}{(1 + \epsilon \cos \phi)^2} d\phi = \frac{1 - \epsilon^2}{q^2} \int_0^{2\pi} r^2 d\phi.$$

This form of expression shows that apart from the factor  $\frac{(1 - \epsilon^2)}{q^2}$

the integral represents twice the area of the elliptical orbit, which, as we know, is  $\pi ab$ , where  $a$  and  $b$  are the semi-axes of the ellipse. Substituting the values previously found for the latter and collecting terms, we find that the first quantum condition gives

$$n'h = -2\pi p_\phi + p_\phi \frac{1 - \epsilon^2}{q^2} 2\pi \frac{q^2}{(1 - \epsilon^2)^{3/2}} = -hk + \frac{2\pi p_\phi}{\sqrt{1 - \epsilon^2}},$$

or

$$\frac{4\pi^2 p_\phi^2}{1 - \epsilon^2} = 4\pi^2 e^2 Z \mu a = h^2 (n' + k)^2.$$

If we put

$$n' + k = n,$$

we have

$$a = \frac{h^2 n^2}{4\pi^2 e^2 \mu Z},$$

so that the energy is

$$E = -\frac{e^2 Z}{2a} = -\frac{2\pi^2 e^4 \mu Z^2}{h^2 n^2} = -\frac{R h}{n^2}.$$

Thus by an argument based on Bohr's theory we have obtained the Balmer term with the correct coefficient.

In deducing this we began by introducing two quantum conditions and hence two quantum numbers, the radial quantum number  $n'$  and the azimuthal quantum number  $k$ . As the orbit has only one period, however, it is found unnecessary to use both quantum numbers in finding the energy levels, as the latter involve only the sum of the two quantum numbers. This we call the *principal quantum number*, as in the unperturbed motion it alone determines the positions of the terms.

The significance of the two other quantum numbers only becomes evident if the degeneracy is removed by some perturbation (due to deviations from the Coulomb field, introduction of the relativistic variation of mass, presence of an external field, or some other cause). We can, however, gain an idea of the meaning of the quantum numbers from the purely geometrical point of view by considering the elliptic orbit. If, as in § 1 (p. 99), we denote the radius of the first circular Bohr orbit for  $Z = 1$  by

$$a_0 = \frac{h^2}{4\pi^2 \mu e^2},$$

the major semi-axis of the ellipse is

$$a = \frac{a_0}{Z} n^2.$$

Using the formulæ above, we similarly obtain the values

$$b = \frac{a_0}{Z} n k, \quad q = \frac{a_0}{Z} k^2$$

for the minor semi-axis and the semi-latus-rectum. The ratio of the axes,  $b/a$ , is accordingly equal to  $k/n$ . If  $n = k$ , we obtain the circular orbits of the Bohr atom, for  $k = 0$  the so-called pendulum orbits (straight lines through the nucleus); but, as was emphasized in the text (p. 108), the latter are to be excluded.

### XV. The Oscillator according to Matrix Mechanics (p. 118).

We proceed to illustrate the fundamental ideas of matrix mechanics by means of an example, namely, the linear harmonic oscillator. We start from the classical expression for the energy,

$$W = \frac{1}{2} \frac{p^2}{\mu} + \frac{1}{2} f q^2,$$

which leads to the equations of motion

$$\dot{p} = -fq, \quad \dot{q} = \frac{p}{\mu}, \quad \text{or} \quad \ddot{q} = -\frac{f}{\mu} q = -\omega_0^2 q, \quad \left( \omega_0 = \sqrt{\frac{f}{\mu}} \right).$$

The difference between classical mechanics and quantum mechanics, as was explained in the text, lies in the fact that the quantities  $p$  and  $q$  are no longer regarded as ordinary functions of the time, but stand for matrices, the element  $q_{nm}$  of which denotes the quantum amplitude associated with the transition from one energy-level  $E_n$  to another,  $E_m$ . Its square, just like  $q^2$ , the square of the amplitude in classical mechanics, is a measure of the intensity of the line of the spectrum emitted in this transition. When we introduce the matrices into the classical equations of motion, we must also bring in the commutation law

$$pq - qp = \frac{h}{2\pi i}$$

as quantum condition.

These equations are easily solved. If the matrix  $(q_{nm})$  is to satisfy the equation of motion  $\ddot{q} + \omega_0^2 q = 0$ , this equation must be satisfied by each element of the matrix independently:

$$\ddot{q}_{nm} + \omega_0^2 q_{nm} = 0.$$

Hence if we make the natural assumption

$$q_{nm} = q_{nm}^{(0)} e^{i\omega_{nm}t},$$

we must have

$$(\omega_0^2 - \omega_{nm}^2) q_{nm} = 0;$$

that is, either  $q_{nm} = 0$  or  $\omega_{nm} = \pm \omega_0$ . Thus all the quantities  $q_{nm}$  vanish except those for which  $\omega_{nm} = +\omega_0$  or  $\omega_{nm} = -\omega_0$ . As we are still free to arrange the numbering of the matrix elements as we please, we make the following statement: the frequency  $\omega_{nm} = +\omega_0$  corresponds to the passage from the  $n$ -th state to the  $(n-1)$ -th (emission),

$\omega_{nm} = -\omega_0$  to the passage from the  $n$ -th state to the  $(n+1)$ -th (absorption), so that

$$\begin{aligned} q_{nm} &= 0 \quad \text{if } m \neq n \pm 1, \\ q_{nm} &\neq 0 \quad \text{if } m = n \pm 1. \end{aligned}$$

This choice of order will be found convenient later, as individual rows (or columns) of the matrices with higher suffixes correspond to states of higher energy. We would, however, expressly emphasize the fact that this fixing of the order does not destroy the generality of the solution in any way.

The co-ordinate matrix accordingly has the form

$$(q_{nm}) = \begin{bmatrix} 0 & q_{01} & 0 & 0 & \dots \\ q_{10} & 0 & q_{12} & 0 & \dots \\ 0 & q_{21} & 0 & q_{23} & \dots \\ \dots & \dots & \dots & \dots & \dots \end{bmatrix}.$$

The momentum matrix has an analogous form; from  $p = \mu\dot{q}$  it follows that

$$p_{nm} = i\mu\omega_{nm}q_{nm},$$

and hence that

$$\begin{aligned} (p_{nm}) &= i\mu \begin{bmatrix} 0 & \omega_{01}q_{01} & 0 & 0 & \dots \\ \omega_{10}q_{10} & 0 & \omega_{12}q_{12} & 0 & \dots \\ 0 & \omega_{21}q_{21} & 0 & \omega_{23}q_{23} & \dots \\ \dots & \dots & \dots & \dots & \dots \end{bmatrix} \\ &= i\mu\omega_0 \begin{bmatrix} 0 & -q_{01} & 0 & 0 & \dots \\ q_{10} & 0 & -q_{12} & 0 & \dots \\ 0 & q_{21} & 0 & -q_{23} & \dots \\ \dots & \dots & \dots & \dots & \dots \end{bmatrix}, \end{aligned}$$

in virtue of the above statement about the quantities  $\omega_{nm}$ .

What we are chiefly interested in is the question of the energy levels. We accordingly calculate the energy matrix, using the co-ordinate matrix and the momentum matrix just given, from the classical energy function

$$W = \frac{1}{2}\frac{p^2}{\mu} + \frac{1}{2}fq^2 = \frac{1}{2\mu}(p^2 + \mu^2\omega_0^2q^2).$$

To do this we have to form the squares  $q^2$  and  $p^2$ . This is carried out according to the rule given in the text (p. 117) for the multiplication of matrices. If  $a$  and  $b$  are two matrices, the elements of their product  $c = ab$  are given by

$$c_{nm} = \sum_k a_{nk}b_{km}.$$

Using this rule, we readily obtain the two expressions

$$q^2 = \begin{bmatrix} q_{01}q_{10} & 0 & q_{01}q_{12} & \dots \\ 0 & q_{10}q_{01} + q_{12}q_{21} & 0 & \dots \\ q_{21}q_{10} & 0 & q_{21}q_{12} + q_{23}q_{32} & \dots \\ \dots & \dots & \dots & \dots \end{bmatrix},$$

$$p^2 = \mu^2\omega_0^2 \begin{bmatrix} q_{01}q_{10} & 0 & -q_{01}q_{12} & \dots \\ 0 & q_{10}q_{01} + q_{12}q_{21} & 0 & \dots \\ -q_{21}q_{10} & 0 & q_{21}q_{12} + q_{23}q_{32} & \dots \\ \dots & \dots & \dots & \dots \end{bmatrix}.$$

Substituting these matrices in the expression for the energy, we obtain the energy matrix

$$W = \mu\omega_0^2 \begin{bmatrix} q_{01}q_{10} & 0 & 0 & \dots \\ 0 & q_{10}q_{01} + q_{12}q_{21} & 0 & \dots \\ 0 & 0 & q_{21}q_{12} + q_{23}q_{32} & \dots \\ \dots & \dots & \dots & \dots \end{bmatrix}.$$

The first thing which strikes us about this matrix is that all the elements vanish except those in the principal diagonal. This, however, is equivalent in meaning to the energy theorem, namely, the energy of a given state is constant as time goes on. Indeed, the time factors disappear from all the diagonal terms; for example, the first term would have the time factor  $e^{i\omega_{01}t}e^{i\omega_{10}t}$ , but this is equal to unity, as  $-\omega_{01} = \omega_{10}$ .

The individual terms of the principal diagonal represent the energies of the individual states, and in fact the element  $W_{nn}$  gives the energy of the  $n$ -th state. The energy matrix involves in the first place the elements of the not yet completely determined co-ordinate matrix. These may, however, be obtained by means of the commutation law  $pq - qp = \hbar/2\pi i$ , and the energy levels may then be found.

If, following the above rules, we substitute the co-ordinate matrix and the momentum matrix in the commutation law, we readily obtain the matrix

$$(pq - qp) = -2i\mu\omega_0 \begin{bmatrix} q_{01}q_{10} & 0 & 0 & \dots \\ 0 & q_{12}q_{21} - q_{10}q_{01} & 0 & \dots \\ 0 & 0 & q_{23}q_{32} - q_{21}q_{12} & \dots \\ \dots & \dots & \dots & \dots \end{bmatrix}.$$

This matrix must be equal to  $\hbar/2\pi i$ , or, more accurately, equal to  $\hbar/2\pi i$  times the unit matrix

$$\begin{bmatrix} 1 & 0 & 0 & \dots \\ 0 & 1 & 0 & \dots \\ 0 & 0 & 1 & \dots \\ \cdot & \cdot & \cdot & \cdot \end{bmatrix}.$$

The two matrices must accordingly have their terms identical; this leads to the system of equations

$$q_{01}q_{10} = \frac{\hbar}{4\pi\mu\omega_0},$$

$$q_{12}q_{21} - q_{01}q_{10} = \frac{\hbar}{4\pi\mu\omega_0},$$

$$q_{23}q_{32} - q_{21}q_{12} = \frac{\hbar}{4\pi\mu\omega_0},$$

$$\cdot \quad \cdot \quad \cdot \quad \cdot \quad \cdot \quad \cdot \quad \cdot \quad \cdot$$

These equations can be solved in succession, giving

$$q_{n,n+1} q_{n+1,n} = (n+1) \frac{\hbar}{4\pi\mu\omega_0}.$$

If we substitute these expressions in the energy matrix, the general diagonal term becomes

$$\begin{aligned} W_{nn} &= E_n = \mu\omega_0^2 (q_{n,n+1} q_{n+1,n} + q_{n,n-1} q_{n-1,n}) \\ &= \frac{\hbar\omega_0}{4\pi} (2n+1) = \hbar\nu_0(n + \tfrac{1}{2}) \quad (n = 0, 1, 2, \dots). \end{aligned}$$

Here again, therefore, we obtain for our term scheme an equidistant succession of energy levels, as in Bohr's theory. The sole difference lies in the fact that the whole term diagram of quantum mechanics is displaced relative to that of Bohr's theory by half a quantum of energy. Although this difference does not manifest itself in the spectrum, it plays a part in statistical problems. In any case it is important to note that the linear harmonic oscillator possesses energy  $\frac{1}{2}\hbar\nu_0$  in the lowest state, the so-called *zero-point energy*.



**XVI. The Oscillator according to Wave Mechanics** (p. 124).

In this section we shall obtain the solution of the wave equation for the linear harmonic oscillator. The equation is

$$\left\{ \frac{d^2}{dq^2} + \frac{8\pi^2 m}{h^2} (E - \tfrac{1}{2} f q^2) \right\} \psi = 0.$$

If we put

$$\lambda = \frac{8\pi^2 m E}{h^2}, \quad \alpha = \frac{2\pi m \omega_0}{h} = \frac{4\pi^2 m \nu_0}{h}, \quad (\omega_0 = 2\pi \nu_0 = \sqrt{f/m})$$

the equation becomes

$$\left\{ \frac{d^2}{dq^2} + \lambda - \alpha^2 q^2 \right\} \psi = 0.$$

Here  $\lambda$  is the proper-value parameter and we have to find the values of  $\lambda$  for which the equation has a finite and unique solution throughout all space.

We can write down one solution of the equation at once, namely, the so-called Gaussian error function

$$\psi = a_0 e^{-\frac{1}{2}\alpha q^2} = a_0 e^{-2\pi^2 m \nu_0 q^2 / h},$$

where the proper-value parameter is  $\lambda_0 = \alpha$ , as we readily see by substitution in the wave equation. The corresponding energy is then given by

$$E_0 = \frac{h^2 \alpha}{8\pi^2 m} = \tfrac{1}{2} h \nu_0.$$

This, we may state in anticipation, is the term of lowest energy, corresponding to the ground state. Here, therefore, we obtain the result given by the matrix method, namely, the ground state possesses a zero-point energy amounting to half Planck's quantum of energy.

In order to find the remaining solutions of the wave equation, it is convenient to assume that  $\psi$  is of the form

$$\psi = e^{-\frac{1}{2}\alpha q^2} v(q).$$

(From the purely mathematical point of view one might be tempted to assume a similar expression but with the plus sign in the exponent; from the physical point of view, however, such an assumption is not permissible, as in that case the wave function would increase beyond all bounds as  $x$  increased.)

If we substitute the above expression in the differential equation, a brief calculation gives the following differential equation for  $v$ :

$$\frac{d^2v}{dq^2} - 2\alpha q \frac{dv}{dq} + (\lambda - \alpha)v = 0.$$

As the equation is homogeneous, the factor  $e^{-\frac{1}{2}\alpha q^2}$  disappears. This differential equation has the following advantage over the wave equation. If we assume that  $v$  may be expressed as a power series in  $q$ ,

$$v = \sum_{\nu=0}^{\infty} a_{\nu} q^{\nu},$$

substitute this series in the differential equation, and equate the coefficients of powers of  $q$  to zero, we obtain the recurrence relation

$$(\nu + 2)(\nu + 1)a_{\nu+2} + \{\lambda - \alpha(2\nu + 1)\}a_{\nu} = 0,$$

involving two coefficients only; whereas if we attempt to satisfy the wave equation by a power series for  $\psi$ , we are led to a recurrence relation involving three coefficients. The recurrence relation just given enables us to calculate the coefficients of the power series.

On detailed examination, however, we find that in general the power series obtained in this way diverges more rapidly than  $e^{\frac{1}{2}\alpha q^2}$  when  $q \rightarrow \infty$ , so that in this case the wave functions would increase beyond all bounds at infinity. The wave functions do not converge, i.e. the power series does not increase more slowly than  $e^{\frac{1}{2}\alpha q^2}$  when  $q \rightarrow \infty$ , unless the power series terminates. This happens for certain special values of the parameter  $\lambda$ ; in fact, as we readily see from the recurrence relation, the series breaks off at the term  $\nu = n$  if the factor multiplying  $a_n$  vanishes, i.e. if

$$\lambda = \lambda_n = \alpha(2n + 1).$$

Recalling the meaning of  $\lambda$ , in this case we obtain the expression

$$E_n = h\nu_0(n + \frac{1}{2})$$

for the energy. It is possible to give a rigorous proof that these values for  $E$  are the only values for which the differential equation has a solution of the required type. The actual solutions take the form of a product of the exponential factor stated above and a finite polynomial in  $q$  of order  $n$ ; in mathematical literature these polynomials are known as Hermite's polynomials. For our purposes it is important to note that wave mechanics and matrix mechanics both give the same energy levels for the harmonic oscillator, namely, a succession of equi-

distant terms, the characteristic deviation from Bohr's theory being the displacement of the terms by  $\frac{1}{2}h\nu_0$ , half the interval between them.

### XVII. The Vibrations of a Circular Membrane (p. 125).

The solution of the differential equation for the vibrations of a circular membrane,

$$\Delta\psi + \lambda\psi = 0 \quad \left( \Delta = \frac{\partial^2}{\partial x^2} + \frac{\partial^2}{\partial y^2} \right),$$

is easily obtained if we use polar co-ordinates. In this case, as we know, the differential equation takes the form

$$\frac{\partial^2\psi}{\partial r^2} + \frac{1}{r} \frac{\partial\psi}{\partial r} + \frac{1}{r^2} \frac{\partial^2\psi}{\partial\phi^2} + \lambda\psi = 0.$$

We see that the solution may be taken as the product of a function  $R$  depending on  $r$  only and a function  $\Phi$  depending on  $\phi$  only; i.e. in polar co-ordinates the variables are separable. The differential equation can then be split up into two equations with the single independent variables  $r$ ,  $\phi$  respectively, by means of a separation parameter, which we shall call  $m^2$ :

$$\frac{d^2R}{dr^2} + \frac{1}{r} \frac{dR}{dr} + \lambda R - \frac{m^2}{r^2} R = 0,$$

$$\frac{d^2\Phi}{d\phi^2} + m^2\Phi = 0.$$

The second equation gives

$$\Phi = \begin{Bmatrix} \cos \\ \sin \end{Bmatrix} (\pm m\phi).$$

It is essential that the wave function shall be one-valued; but this is not the case unless  $m$  is an integer, for otherwise an increase of  $2\pi$  in the value of  $\phi$  would give a different value for the wave function; hence the separation parameter  $m$  must be an integer ( $m \geq 0$ ).

The differential equation in  $r$  is the well-known equation defining the Bessel function  $J_m(\sqrt{\lambda}r)$  of the  $m$ -th order with argument  $\sqrt{\lambda}r$ . Here, however, we must take the boundary conditions into account. As a particular case, we assume that the membrane is fixed at the circumference, so that for all points on the boundary  $R(\rho) = 0$ , where  $\rho$  is the radius of the membrane. Now the Bessel function of any order has an infinite number of zeros (in fact if the value of the argument

is not too small its graph is similar to a sine wave); the zeros of the Bessel function  $J_m(z)$  we shall call  $z_0, z_1, z_2, \dots$ . In order that the boundary condition may be satisfied, the argument of the Bessel functions must coincide with one of these zeros when  $r = \rho$ . Hence not all values of the parameter are possible, but only those for which

$$\sqrt{\lambda} \rho = z_0, \text{ or } z_1, \text{ or } z_2, \dots$$

In the first case the wave function has no zeros other than those for  $r = \rho$  and (where  $m > 0$ ) for  $r = 0$ ; in the second case it has one other zero, in the third case two other zeros; hence in general for the proper value

$$\sqrt{\lambda_n} \rho = z_n$$

there are  $n$  other zeros (circular nodal lines; see fig. 20, p. 125).

### XVIII. Solution of Schrödinger's Equation for the Kepler Problem (p. 125).

Schrödinger's equation for the Kepler problem is

$$\Delta\psi + \frac{8\pi^2m}{h^2} \left( E + \frac{e^2Z}{r} \right) \psi = 0.$$

Changing to three-dimensional polar co-ordinates  $r, \theta, \phi$ , we obtain the equation

$$\left\{ \frac{\partial^2}{\partial r^2} + \frac{2}{r} \frac{\partial}{\partial r} + \frac{1}{r^2} \frac{\partial^2}{\partial \theta^2} + \frac{\cot \theta}{r^2} \frac{\partial}{\partial \theta} + \frac{1}{r^2 \sin^2 \theta} \frac{\partial^2}{\partial \phi^2} \right\} \psi + \frac{8\pi^2m}{h^2} \left( E + \frac{e^2Z}{r} \right) \psi = 0.$$

This equation is likewise separable; if we put

$$\psi = R(r)\Theta(\theta)\Phi(\phi),$$

the equation may be split up into the three equations

$$\left\{ \frac{d^2}{dr^2} + \frac{2}{r} \frac{d}{dr} + \frac{8\pi^2m}{h^2} \left( E + \frac{e^2Z}{r} \right) - \frac{\lambda}{r^2} \right\} R = 0,$$

$$\left\{ \frac{d^2}{d\theta^2} + \cot \theta \frac{d}{d\theta} + \lambda - \frac{m^2}{\sin^2 \theta} \right\} \Theta = 0,$$

$$\left\{ \frac{d^2}{d\phi^2} + m^2 \right\} \Phi = 0,$$

where  $m$  (as before) and  $\lambda$  are the separation parameters. The solution of the third equation we know already:

$$\Phi = \begin{cases} \cos \\ \sin \end{cases} (m\phi), \text{ or } \Phi = e^{im\phi},$$

where  $m$  must be an integer, otherwise  $\Phi$  is not one-valued.

The second equation is the equation defining the spherical harmonic  $P_l^m(\cos\theta)$ , when  $\lambda$  has the value  $l(l+1)$  and  $|m| \leq l$ ; for other values the equation has no finite one-valued solution. We shall prove this generally by combining the  $\theta$ - and  $\phi$ -relationships in the general surface harmonic  $Y_l(\theta, \phi)$ . If for brevity we introduce the notation

$$\Lambda = \frac{1}{\sin\theta} \frac{\partial}{\partial\theta} \sin\theta \frac{\partial}{\partial\theta} + \frac{1}{\sin^2\theta} \frac{\partial^2}{\partial\phi^2},$$

so that

$$\Delta = \frac{\partial^2}{\partial r^2} + \frac{2}{r} \frac{\partial}{\partial r} + \frac{\Lambda}{r^2},$$

$Y_l(\theta, \phi)$  must satisfy the differential equation

$$\Lambda Y_l + \lambda Y_l = 0.$$

A general solution of this equation is obtained in the following way. We consider a homogeneous polynomial  $U_l$  of the  $l$ -th degree in  $x, y, z$ , which satisfies Laplace's equation

$$\Delta U_l = 0.$$

If we now define a function  $Y_l$  of the ratios  $x/r, y/r, z/r$ , i.e. of the angles  $\theta$  and  $\phi$  only, by means of the equation

$$U_l = r^l Y_l,$$

substitute this in Laplace's equation  $\Delta U_l = 0$ , and carry out the differentiation with respect to  $r$ , we obtain the equation

$$\left( \frac{\partial^2}{\partial r^2} + \frac{2}{r} \frac{\partial}{\partial r} \right) r^l Y_l + \frac{\Lambda}{r^2} r^l Y_l = r^{l-2} \{ \Lambda Y_l + l(l+1) Y_l \} = 0;$$

that is, the functions defined above are solutions of the differential equation  $\Delta Y_l + \lambda Y_l = 0$  in the case where

$$\lambda = l(l+1).$$

It is possible to prove that no other values of  $\lambda$  give finite and continuous one-valued solutions of the differential equation. The proper values of the equation

$$\Lambda Y_l + \lambda Y_l = 0$$

are accordingly equal to  $l(l+1)$ . Moreover, it is easy to reckon up the number of arbitrary parameters in any function of the  $l$ -th degree. The most general homogeneous polynomial of the  $l$ -th degree in  $x, y, z$  contains  $\frac{1}{2}l(l+1)(l+2)$  arbitrary constants (it contains one term in  $x^l$ , two in  $x^{l-1}$ , three in  $x^{l-2}$ , and so on, and finally  $(l+1)$  terms not containing  $x$  at all). These constants, however, are connected by certain relations depending on the condition  $\Delta U_l = 0$ ; this equation is equivalent to  $\frac{1}{2}l(l-1)$  equations for determining the coefficients, for  $\Delta U_l$  is a homogeneous function of the  $(l-2)$ -th degree, which must vanish identically. Hence  $U_l$  contains

$$\frac{1}{2}\{l(l+1)(l+2) - l(l-1)\} = 2l+1$$

independent coefficients. Accordingly there are  $(2l+1)$  linearly independent spherical harmonics of the  $l$ -th degree. If we write them in the usual form  $Y_l^{(m)} = P_l^m e^{im\phi}$ , they correspond to the  $(2l+1)$  possible values of the third (magnetic) quantum number  $m$ .

We now pass on to the differential equation for the radial function  $R$ :

$$\left\{ \frac{d^2}{dr^2} + \frac{2}{r} \frac{d}{dr} + \frac{8\pi^2 m}{h^2} \left( E + \frac{e^2 Z}{r} \right) - \frac{l(l+1)}{r^2} \right\} R = 0.$$

Its solution must be finite and continuous for all values of  $r$  from zero to infinity. Here we are chiefly interested in the magnitude of the proper values  $E$  for which this equation has a solution satisfying the prescribed conditions. In particular, we shall only discuss the case where  $E < 0$ . This corresponds to the elliptical orbits in Bohr's theory; energy must be supplied in order to remove the electron to the boundary of the atom or, better expressed, to an infinite distance from the nucleus. The case where  $E > 0$  would correspond to Bohr's hyperbolic orbits.

For the sake of simplicity we shall introduce rational units. We shall measure the radius in multiples of the Bohr radius  $\frac{h^2}{4\pi^2 m e^2 Z}$  and the energy in multiples of the ground state of the Bohr atom,  $-\frac{2\pi^2 m e^4 Z^2}{h^2}$ , i.e. we put (p. 99)

$$r = \rho \frac{h^2}{4\pi^2 m e^2 Z}, \quad E = \epsilon \left( \frac{-2\pi^2 m e^4 Z^2}{h^2} \right).$$

The wave equation then takes the simpler form

$$\left\{ \frac{d^2}{d\rho^2} + \frac{2}{\rho} \frac{d}{d\rho} - \epsilon + \frac{2}{\rho} - \frac{l(l+1)}{\rho^2} \right\} R = 0.$$

We begin by investigating the behaviour of the function  $R$  for very large values of  $\rho$ . We accordingly omit the terms in  $1/\rho$  and  $1/\rho^2$  in the differential equation, so that the behaviour at infinity is given by the equation  $\left\{ \frac{d^2}{d\rho^2} - \epsilon \right\} R_\infty = 0$ . Its solution is  $R_\infty = e^{\pm \rho \sqrt{\epsilon}}$ ; the upper sign, however, is impossible, as then the wave function would increase exponentially beyond all bounds as  $r$  and  $\rho$  increased and therefore could not be regarded as a proper function.

The second special region is that of the origin. We obtain an approximation by omitting terms of the differential equation which tend to infinity more slowly than  $1/\rho^2$  when  $\rho \rightarrow 0$ . The equation is then

$$\left\{ \frac{d^2}{d\rho^2} + \frac{2}{\rho} \frac{d}{d\rho} - \frac{l(l+1)}{\rho^2} \right\} R_0 = 0.$$

Here the solutions are  $R_0 = \rho^l$  and  $R_0 = \rho^{-l-1}$ ; the second of these is impossible, for it becomes infinite at the origin.

We now know how the desired function behaves at the two singularities  $\rho = 0$  and  $\rho = \infty$ . It is natural to assume that the complete function  $R$  is of the form

$$R = e^{-\rho \sqrt{\epsilon}} \rho^l f(\rho),$$

where  $f$  is a function of  $\rho$  which must of course behave regularly at these two points (i.e. must not increase more rapidly than  $e^{+\rho \sqrt{\epsilon}}$  at infinity) and which determines the nature of the complete function in the region intermediate between the regions of validity of the power,  $\rho^l$ , and of the exponential function. Substituting this expression in the differential equation, we readily obtain the following differential equation for  $f$ :

$$\frac{d^2 f}{d\rho^2} + \frac{2(l+1)}{\rho} \frac{df}{d\rho} - 2\sqrt{\epsilon} \frac{df}{d\rho} + \frac{2}{\rho} (1 - \sqrt{\epsilon}(l+1))f = 0.$$

We attempt to solve it by means of a power series in  $\rho$  (or, better, in  $2\rho\sqrt{\epsilon}$ ) and accordingly write

$$f = \sum_0^{\infty} a_\nu (2\rho\sqrt{\epsilon})^\nu.$$

Substituting this in the differential equation and rearranging the terms somewhat, we obtain

$$\sum_{\nu=0}^{\infty} a_\nu (2\rho\sqrt{\epsilon})^{\nu-2} \nu(\nu+2l+1) - \sum_{\nu=0}^{\infty} a_\nu (2\rho\sqrt{\epsilon})^{\nu-1} \left( \nu+l+1 - \frac{1}{\sqrt{\epsilon}} \right) = 0.$$

This series must vanish identically, so that we obtain the recurrence relation

$$a_{\nu+1}(\nu+1)(\nu+2l+2) = a_{\nu} \left( \nu+l+1 - \frac{1}{\sqrt{\epsilon}} \right)$$

for the coefficients. At the origin the function  $f$  is of course finite and equal to the initial term  $a_0$ . At infinity, on the other hand,  $f$  becomes infinite, and in fact, as more detailed analysis shows, to a higher order than  $e^{+\rho\sqrt{\epsilon}}$ , provided the series for  $f$  does not terminate. If, however, the series does terminate,  $R$  vanishes at infinity, despite the fact that  $f$  becomes infinite, in virtue of the exponential factor  $e^{-\rho\sqrt{\epsilon}}$ . The condition that the series shall terminate is obtained from the recurrence relation; the series breaks off after the  $n_r$ -th term if

$$n_r + l + 1 = \frac{1}{\sqrt{\epsilon}}.$$

That is,  $1/\sqrt{\epsilon}$  must be a positive integer, or

$$\epsilon = \frac{1}{n^2},$$

where  $n = n_r + l + 1$ ;  $n$  is known as the principal quantum number,  $n_r$  as the radial quantum number.

We see, therefore, that solutions of the differential equation which satisfy the conditions of finiteness, continuity, and one-valuedness can be found only for certain discrete values of the parameter  $\epsilon$ , namely the values  $\epsilon = 1/n^2$ . Hence certain definite energy levels are alone possible, namely

$$E = -\frac{2\pi^2 me^4 Z^2}{h^2 n^2} = -\frac{h R_0 Z^2}{n^2},$$

which are those introduced in Bohr's theory to explain spectra.

We may add that the polynomials  $f$  are known as Laguerre's polynomials; these, however, we shall not discuss in more detail here, except to mention that their zeros determine the position of the nodal surfaces  $r = \text{const.}$ ; in fact,  $R$  has  $n_r$  nodes, not counting the zeros for  $r = 0$  (in the case where  $l > 0$ ) and  $r = \infty$ .

### XIX. The Resultant Angular Momentum (p. 129).

In the text we gave the operators

$$M_x = \frac{h}{2\pi i} \left( y \frac{\partial}{\partial z} - z \frac{\partial}{\partial y} \right)$$

. . . . .



corresponding to the components of the angular momentum. The square of the resultant angular momentum is

$$\mathbf{M}^2 = \mathbf{M}_x^2 + \mathbf{M}_y^2 + \mathbf{M}_z^2,$$

where the expression  $\mathbf{M}_x^2$  implies the repetition of the operator  $\mathbf{M}_x$ ; that is, if we have operated with  $\mathbf{M}_x$  on a function  $\psi$ , the application of  $\mathbf{M}_x^2$  to the function  $\psi$  means the formation of  $\mathbf{M}_x(\mathbf{M}_x\psi)$ .

For the first component, for example, we accordingly have

$$\mathbf{M}_x^2\psi = -\left(\frac{h}{2\pi}\right)^2 \left(y \frac{\partial}{\partial z} - z \frac{\partial}{\partial y}\right) \left(y \frac{\partial}{\partial z} - z \frac{\partial}{\partial y}\right) \psi,$$

or, carrying out the differentiations,

$$\mathbf{M}_x^2\psi = \left(\frac{h}{2\pi}\right)^2 \left(-y^2 \frac{\partial^2 \psi}{\partial z^2} + 2yz \frac{\partial}{\partial y} \frac{\partial \psi}{\partial z} - z^2 \frac{\partial^2 \psi}{\partial y^2} + y \frac{\partial \psi}{\partial y} + z \frac{\partial \psi}{\partial z}\right);$$

hence the operator  $\mathbf{M}_x^2$  is identical with the operator

$$\left(\frac{h}{2\pi}\right)^2 \left(-y^2 \frac{\partial^2}{\partial z^2} - z^2 \frac{\partial^2}{\partial y^2} + 2yz \frac{\partial}{\partial y} \frac{\partial}{\partial z} + y \frac{\partial}{\partial y} + z \frac{\partial}{\partial z}\right).$$

Adding the three components, we have

$$\begin{aligned} \mathbf{M}^2 = \left(\frac{h}{2\pi}\right)^2 & \left\{ -\left[(y^2 + z^2) \frac{\partial^2}{\partial x^2} + (z^2 + x^2) \frac{\partial^2}{\partial y^2} + (x^2 + y^2) \frac{\partial^2}{\partial z^2}\right] \right. \\ & + 2\left[xy \frac{\partial}{\partial x} \frac{\partial}{\partial y} + yz \frac{\partial}{\partial y} \frac{\partial}{\partial z} + zx \frac{\partial}{\partial z} \frac{\partial}{\partial x}\right] \\ & \left. + 2\left[x \frac{\partial}{\partial x} + y \frac{\partial}{\partial y} + z \frac{\partial}{\partial z}\right] \right\}. \end{aligned}$$

or, rearranging the terms and noting that

$$\left(x \frac{\partial}{\partial x}\right)^2 = x \frac{\partial}{\partial x} x \frac{\partial}{\partial x} = x^2 \frac{\partial^2}{\partial x^2} + x \frac{\partial}{\partial x},$$

$$\begin{aligned} \mathbf{M}^2 = \left(\frac{h}{2\pi}\right)^2 & \left\{ -r^2 \left( \frac{\partial^2}{\partial x^2} + \frac{\partial^2}{\partial y^2} + \frac{\partial^2}{\partial z^2} \right) + \left( x \frac{\partial}{\partial x} + y \frac{\partial}{\partial y} + z \frac{\partial}{\partial z} \right)^2 \right. \\ & \left. + \left( x \frac{\partial}{\partial x} + y \frac{\partial}{\partial y} + z \frac{\partial}{\partial z} \right) \right\}. \end{aligned}$$

As we know,  $x \frac{\partial}{\partial x} + y \frac{\partial}{\partial y} + z \frac{\partial}{\partial z} = \frac{\partial}{\partial r}$ , so that

$$M^2 = \left(\frac{h}{2\pi}\right)^2 \left\{ -r^2 \Delta + r \frac{\partial}{\partial r} r \frac{\partial}{\partial r} + r \frac{\partial}{\partial r} \right\}.$$

If we now change to polar co-ordinates, we see (cf. Appendix XVIII, p. 298) that

$$\begin{aligned} M^2 &= \left(\frac{h}{2\pi}\right)^2 \left\{ -r^2 \left( \frac{\partial^2}{\partial r^2} + \frac{2}{r} \frac{\partial}{\partial r} \right) - \Lambda + r \frac{\partial}{\partial r} r \frac{\partial}{\partial r} + r \frac{\partial}{\partial r} \right\} \\ &= -\frac{h^2}{4\pi^2} \Lambda. \end{aligned}$$

This simply means that for every state with azimuthal quantum number  $l$ ,  $M^2$  has the proper value  $\frac{h^2}{4\pi^2} l(l+1)$ . The magnitude of the resultant angular momentum is accordingly quantised and has the value  $\frac{h}{2\pi} \sqrt{l(l+1)}$ .

Further, in polar co-ordinates with polar axis  $z$  we have

$$M_z = \frac{h}{2\pi i} \left( x \frac{\partial}{\partial y} - y \frac{\partial}{\partial x} \right) = \frac{h}{2\pi i} \frac{\partial}{\partial \phi}.$$

Applying this operator to the one-valued function  $e^{im\phi}$  (where  $m = 0, \pm 1, \pm 2, \&c.$ ), we have

$$M_z e^{im\phi} = \frac{h}{2\pi} m e^{im\phi},$$

that is, the component  $M_z$  of the angular momentum is also quantised, its proper values being integral multiples of  $h/2\pi$ , Bohr's unit of angular momentum. We cannot go into the behaviour of the other two components ( $M_x, M_y$ ) here, but merely mention that the matrices corresponding to them can be evaluated by the method explained in Appendix XXI.

## XX. Deduction of Rutherford's Scattering Formula by Wave Mechanics (p. 131).

We consider Schrödinger's differential equation

$$\left\{ -\frac{h^2}{8\pi^2 m} \Delta + V + \frac{h}{2\pi i} \frac{\partial}{\partial t} \right\} \psi = 0,$$

and with it the conjugate imaginary form

$$\left\{ -\frac{h^2}{8\pi^2 m} \Delta + V - \frac{h}{2\pi i} \frac{\partial}{\partial t} \right\} \psi^* = 0.$$

If we multiply the first equation by  $\psi^*$ , the second by  $\psi$ , and subtract, we obtain

$$-\frac{\hbar^2}{8\pi^2m}(\psi^*\Delta\psi - \psi\Delta\psi^*) + \frac{\hbar}{2\pi i}\left(\psi^*\frac{\partial\psi}{\partial t} + \psi\frac{\partial\psi^*}{\partial t}\right) = 0,$$

or

$$\frac{\partial}{\partial t}(\psi\psi^*) = \frac{\hbar i}{4\pi m}\left\{\frac{\partial}{\partial x}\left(\psi^*\frac{\partial\psi}{\partial x} - \psi\frac{\partial\psi^*}{\partial x}\right) + \dots\right\}.$$

Thus, defining the stream vector  $\mathbf{s}$  (number of particles per square centimetre per second) by the components

$$s_x = \frac{\hbar}{2\pi i} \frac{1}{2m} \left( \psi^* \frac{\partial\psi}{\partial x} - \psi \frac{\partial\psi^*}{\partial x} \right), \quad s_y = \dots, \quad s_z = \dots,$$

we have

$$\frac{\partial}{\partial t} |\psi|^2 = -\text{div } \mathbf{s};$$

hence

$$\frac{d}{dt} \int |\psi|^2 dv = - \int \text{div } \mathbf{s} dv = - \int s_n d\sigma.$$

Since the last integral represents the total number of particles leaving the bounding surface per second, this relation is consistent with the assumption that

$$\int |\psi|^2 dv$$

represents the number of particles within the volume  $v$ . If  $\psi$ , and therefore  $\mathbf{s}$ , vanishes at the boundary (for example, at infinity), then

$$\int |\psi|^2 dv = \text{const.},$$

i.e. the total number of particles remains constant.

For a stationary process ( $\psi \sim e^{-(2\pi i/\hbar)Et}$ ), in which particles with large energy come from infinity, we have approximately, as in the text,

$$\left( \frac{\hbar^2}{8\pi^2m} \Delta + E \right) \psi = V e^{(2\pi i/\hbar)pz}$$

or

$$\Delta\psi + k^2\psi = F(x, y, z)e^{ikz},$$

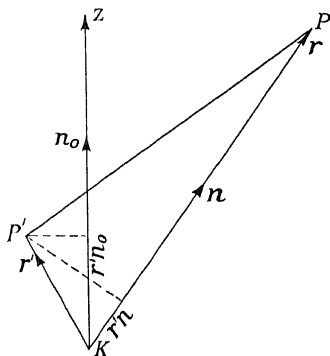


Fig. 5.—Diagram for Rutherford's scattering formula

where

$$k = \frac{2\pi}{h} p = \frac{2\pi}{h} \sqrt{(2mE)}, \quad F = \frac{8\pi^2 m}{h^2} V.$$

The general solution, which corresponds to superposition of the incident wave  $e^{ikz}$  with an outgoing spherical wave, is

$$\psi = e^{ikz} - \frac{1}{4\pi} \iint \frac{e^{ikR}}{R} F(r') e^{ikz'} dx' dy' dz',$$

where  $R$  is the distance between  $P(x, y, z)$  and  $P'(x', y', z')$ .

It is sufficient to have the solution for very great distances  $r = \sqrt{(x^2 + y^2 + z^2)}$  from the nucleus  $K$  (fig. 5); we have then approximately

$$R = r - \mathbf{r}' \cdot \mathbf{n},$$

where  $\mathbf{r}'$  is the vector  $KP'$ , and  $\mathbf{n}$  is the unit vector in the direction  $KP$ . We can also put

$$z' = \mathbf{r}' \cdot \mathbf{n}_0,$$

where  $\mathbf{n}_0$  is the unit vector in the  $z$ -direction. Hence

$$\psi = e^{ikz} - \frac{e^{ikr}}{4\pi r} \iint e^{ik\mathbf{r}'(\mathbf{n}_0 - \mathbf{n})} F(r') dx' dy' dz'.$$

On introducing polar co-ordinates  $\alpha, \beta$  round the vector  $\mathbf{n}_0 - \mathbf{n}$  as axis, we find

$$\psi = e^{ikz} - \frac{e^{ikr}}{r} f(\theta);$$

here

$$f(\theta) = \frac{1}{4\pi} \int_0^{2\pi} d\beta \int_0^\pi \sin \alpha d\alpha \int_0^\infty r'^2 dr' e^{iK\mathbf{r}' \cos \alpha} F(r'),$$

with

$$\begin{aligned} K &= k |\mathbf{n}_0 - \mathbf{n}| = k \sqrt{n_x^2 + n_y^2 + (1 - n_z)^2} \\ &= k \sqrt{2(1 - n_z)} = k \sqrt{2(1 - \cos \theta)} = 2k \sin \frac{\theta}{2}, \end{aligned}$$

where  $\theta$  is the angle between  $\mathbf{n}$ , i.e.  $KP$ , and the  $z$ -axis. The integrations with respect to  $\alpha$  and  $\beta$  can be carried out:

$$f(\theta) = \frac{8\pi^2 m}{h^2} \int_0^\infty r'^2 V(r') dr' \cdot \frac{\sin Kr'}{Kr'}.$$

## APPENDIX XX

The intensity of the incident wave is

$$s_z^0 = \frac{h}{2\pi i} \left( \frac{ik}{m} \right);$$

that of the scattered wave (for  $r$  great) is found from the secondary wave:

$$\begin{aligned} s_r &= \frac{h}{2\pi i} \frac{1}{2m} \left( \psi^* \frac{\partial \psi}{\partial r} - \psi \frac{\partial \psi^*}{\partial r} \right) \\ &= \frac{h}{2\pi i} \frac{\{f(\theta)\}^2}{r^2} \left( \frac{ik}{m} \right). \end{aligned}$$

Hence the scattering probability per unit solid angle is

$$w = r^2 \frac{s_r}{s_z^0} = \{f(\theta)\}^2.$$

The potential can be represented approximately by

$$V = ZeE \frac{e^{-r/a}}{r},$$

where the exponential function represents the screening effect of the bound electrons;  $a$  is of the order of magnitude of atomic radii,  $10^{-8}$  cm. The  $r$ -integration can then also be effected, and we get

$$f(\theta) = \frac{8\pi^2 m}{h^2} \frac{ZeE}{K^2 + 1/a^2}$$

Now  $K = 2k \sin \frac{\theta}{2} = \frac{4\pi}{\lambda} \sin \frac{\theta}{2}$ , where  $\lambda$  is the de Broglie wave-length; hence  $1/a^2$  is always small compared with  $K$  for particles with velocities at all rapid, except for the immediate neighbourhood of the direction of the primary beam ( $\theta \sim 0$ ). If  $1/a^2$  is neglected, the effect of the screening disappears entirely, and we find

$$f(\theta) = \frac{8\pi^2 m}{h^2} \frac{ZeE}{K^2} = \frac{2\pi^2 m ZeE}{h^2 k^2 \sin^2 \theta/2},$$

or, with

$$k = \frac{2\pi}{h} p = \frac{2\pi}{h} mv \quad (v = \text{velocity});$$

$$f(\theta) = \frac{ZeE}{2mv^2} \frac{1}{\sin^2 \theta/2}.$$

On squaring, this gives exactly Rutherford's scattering formula,

p. 53 (apart from the notation:  $m$  instead of  $M$ ,  $\theta$  instead of  $\phi$ ).

Whatever the law for  $V(r)$  may be, a closed formula can be obtained (Bethe) just as easily as in the case before us, this law being expressed by an integral taken over the appropriate screening charge density, the so-called atomic scattering factor.

### XXI. Deduction of the Selection Rules for the Kepler Problem (p. 133).

As was shown in § 6 (p. 130), the radiation associated with a quantum jump is given essentially by the matrix element of the co-ordinate, which is related to the wave-mechanical mean value of the electrical dipole moment in the way indicated in the text:

$$p_x = \int x \psi_{n'l'm'}^* \psi_{nlm} dx dy dz,$$

and so on. Here we shall prove that a large number of these integrals vanish, and that it is only for certain combinations of the quantum numbers  $(l, m)$  and  $(l', m')$  satisfying the "selection rules" that non-zero values are obtained for these integrals.

To do this we must investigate the proper functions for the various states. We deduced these in Appendix XVIII (p. 298), and showed that they may be written in the form

$$\psi_{n,l} = R_{n,l}(r) Y_l(\theta, \phi),$$

where  $R_{n,l}(r)$  is a function of the radius only and  $Y_l(\theta, \phi)$  denotes a general surface harmonic satisfying the differential equation

$$\Delta Y_l + l(l+1)Y_l = 0 \quad \left( \Delta = \frac{1}{\sin \theta} \frac{\partial}{\partial \theta} \sin \theta \frac{\partial}{\partial \theta} + \frac{1}{\sin^2 \theta} \frac{\partial^2}{\partial \phi^2} \right).$$

We recall that the functions  $Y_l$  are obtained by removing the factor  $r^l$  from a homogeneous polynomial of the  $l$ -th degree,  $U_l(x, y, z)$ , which satisfies Laplace's equation  $\Delta U_l = 0$ . If we introduce this form of the wave functions into the matrix element above, the integral is split up into two parts, an integral over the elementary solid angle  $d\omega = \sin \theta d\theta d\phi$ , which is of the form

$$J_{ll'}^{(x)} = \int_{\Omega} Y_{l'}^* Y_l d\omega$$

and therefore does not depend on  $r$ , and the radial integral

$$\int R_{n'l'}^* R_{nl} r^3 dr$$

(for  $dx dy dz = r^2 dr d\omega$ ). Here we shall consider the first integral only, as it gives the selection rules by itself, the second integral merely

determining the gradations in the intensity of the radiation for the various transitions.

We now assert that  $J_{l'l'}^{(\infty)}$  is always zero unless the selection rule  $l' = l \pm 1$  is satisfied. To prove this some preliminaries are required. In the first place, it is easy to see that the integral

$$N_{l'l'} = \int Y_l^* Y_{l'} d\omega$$

is zero except when  $l = l'$ . For the differential equation

$$Y_l^* \Lambda Y_{l'} - Y_{l'} \Lambda Y_l^* = [l'(l' + 1) - l(l + 1)] Y_l^* Y_{l'}$$

when  $\Lambda$  is written out in full becomes

$$\begin{aligned} \frac{1}{\sin \theta} \frac{\partial}{\partial \theta} \left[ \sin \theta \left( Y_l^* \frac{\partial Y_{l'}}{\partial \theta} - Y_{l'} \frac{\partial Y_l^*}{\partial \theta} \right) \right] + \frac{1}{\sin^2 \theta} \frac{\partial}{\partial \phi} \left( Y_l^* \frac{\partial Y_{l'}}{\partial \phi} - Y_{l'} \frac{\partial Y_l^*}{\partial \phi} \right) \\ = [l'(l' + 1) - l(l + 1)] Y_l^* Y_{l'}; \end{aligned}$$

if we now multiply both sides by  $d\omega = \sin \theta d\theta d\phi$  and integrate over the whole range of  $\theta$  and  $\phi$ , the left-hand side vanishes, the first term on account of the vanishing of  $\sin \theta$  at the limits, the second term on account of its periodicity in  $\phi$ . Hence the integral of the right-hand side must vanish also, and the fact that

$$[l'(l' + 1) - l(l + 1)] \int Y_l^* Y_{l'} d\omega = 0$$

leads to the result stated above, since the factor  $l'(l' + 1) - l(l + 1)$  is not zero unless  $l' = l$  ( $l$  and  $l'$  are not less than zero); if  $l' \neq l$ , therefore, the integral must vanish.

The proof of the selection rules is now completed as follows: as we shall show immediately, an expression of the form  $\frac{x}{r} Y_l$  may always be represented as the sum of two general surface harmonics of orders  $(l + 1)$  and  $(l - 1)$ , that is,

$$\frac{x}{r} Y_l = Y_{l+1} + Y_{l-1}.$$

If we substitute this in the integral  $J_{l'l'}^{(\infty)}$ , it follows from the relation just proved that the integral vanishes identically except when  $l' = l \pm 1$ .

The proof of the relation just used follows from the theorem that every homogeneous polynomial  $F(x, y, z)$  of the  $n$ -th degree can be reduced in one way only to an expression of the form

$$F = U_n + r^2 U_{n-2} + r^4 U_{n-4} + \dots + r^{2h} U_{n-2h} + \dots,$$

where the functions  $U_n$  are the potential functions introduced above. If  $n = 0$  or  $1$  the theorem is trivial, as every polynomial of zero or first degree is already a potential function. We shall prove by induction that it is true in general. If we assume that the theorem is true for all polynomials of degree less than  $n$ , it is certain that  $\Delta F$ , which is a polynomial of the  $(n - 2)$ -th degree, can be written in the form

$$\Delta F = U_{n-2}^* + r^2 U_{n-4}^* + \dots + r^{2h-2} U_{n-2h}^* + \dots,$$

where the functions  $U_n^*$  are potential functions different from the  $U_n$ 's. We assert that the general solution of this equation, subject to the condition that  $F$  is of degree  $n$ , is

$$F = G + \frac{r^2 U_{n-2}^*}{2(2n-1)} + \frac{r^4 U_{n-4}^*}{4(2n-3)} + \dots + \frac{r^{2h} U_{n-2h}^*}{2h(2n-2h+1)} + \dots,$$

where  $G$  is an arbitrary function of the  $n$ -th degree, which satisfies Laplace's equation  $\Delta G = 0$ ; for if we pick out an arbitrary term of this series and make use of Euler's theorem for homogeneous polynomials, we have

$$\begin{aligned} \Delta(r^{2h} U_{n-2h}^*) &= r^{2h} \Delta U_{n-2h}^* + 2 \operatorname{grad} r^{2h} \operatorname{grad} U_{n-2h}^* + (\Delta r^{2h}) U_{n-2h}^* \\ &= 0 + 2 \cdot 2hr^{2h-2}(n-2h) U_{n-2h}^* + 2h(2h+1)r^{2h-2} U_{n-2h}^* \\ &= 2h(2n-2h+1)r^{2h-2} U_{n-2h}^*, \end{aligned}$$

which gives the result stated. The polynomial  $F$  is therefore reduced to the required form if we put  $U_{n-2h}^* = 2h(2n-2h+1)U_{n-2h}$  and  $U_n = G$ .

The proof of the relation stated follows at once:  $xU_l$  is a polynomial of the  $(l+1)$ -th degree and may therefore be expressed in the form given above: as  $\Delta(xU_l) = 2 \frac{\partial U_l}{\partial x}$  is itself a potential function (for  $\Delta \frac{\partial U_l}{\partial x} = \frac{\partial}{\partial x} \Delta U_l$ ) of degree  $l-1$ , only the first term of the series, namely  $U_{l-1}$ , appears in it, so that  $xU_l$  involves only the first two terms of the series: hence

$$xU_l = U_{l+1} + r^2 U_{l-1},$$

or, using the functions  $Y_l$  and dividing the equation by  $r^{l+1}$ ,

$$\frac{x}{r} Y_l = Y_{l+1} + Y_{l-1},$$

as we asserted above.



Now that we have given a general proof of the validity of the selection rule  $l' = l \pm 1$  for the azimuthal quantum number, we shall go on to consider the case where the levels with differing magnetic quantum numbers are split up, e.g. by an external magnetic field, and seek to find the selection rules which apply to transitions between these various levels.

For this purpose we change to the usual notation for the wave functions, in which the magnitude of the angular momentum about a particular axis may be recognized directly, viz. the general surface harmonic  $Y_l(\theta, \phi)$  is replaced by the special function  $P_l^m e^{im\phi}$ ; then the angular momentum about the polar axis is equal to  $m\hbar/2\pi$ . The integral  $J_{l'l'm'm'}$  then splits up into two integrals involving  $\theta$  alone and  $\phi$  alone; here we are interested only in the integral in  $\phi$ , the integral in  $\theta$  merely giving the selection rule just obtained for  $l$ .

We accordingly have to consider the three integrals

$$I_x = \int_0^{2\pi} x e^{i(m-m')\phi} d\phi,$$

$$I_y = \int_0^{2\pi} y e^{i(m-m')\phi} d\phi,$$

$$I_z = \int_0^{2\pi} z e^{i(m-m')\phi} d\phi,$$

which are readily evaluated if we introduce polar co-ordinates instead of  $x, y$ , and  $z$ . Since  $z = r \cos \theta$ , the third integral is zero except when  $m' = m$ . We combine the two others in the following way:

$$\begin{aligned} I_x \pm iI_y &= \int_0^{2\pi} (x \pm iy) e^{i(m-m')\phi} d\phi = r \sin \theta \int_0^{2\pi} (\cos \phi \pm i \sin \phi) e^{i(m-m')\phi} d\phi, \\ &= r \sin \theta \int_0^{2\pi} e^{i(m-m' \pm 1)\phi} d\phi. \end{aligned}$$

This gives the selection rule  $m' = m \pm 1$ .

What, then, is the meaning of these two different selection rules in the case of the matrix elements  $\bar{x}$ ,  $\bar{y}$ ,  $\bar{z}$  for the emitted radiation? We saw above that apart from the factor  $e$  (the elementary charge) this matrix element means the wave-mechanical average (probability) of the electrical dipole moment for a transition from the state  $n\bar{l}m$  to the state  $n'\bar{l}'m'$ . According to the correspondence principle, however, the existence of a dipole moment variable in time implies the emission of electromagnetic radiation; also, the radiation field of an oscillating dipole exhibits the peculiarity that there is no emission in

the direction of its vibration. If, therefore, the selection rule  $m' = m$  holds only for the  $z$ -component of the dipole moment, this means that the radiation corresponding to the transition  $m \rightarrow m$  only appears when the electrical charge-cloud vibrates in the special direction (the  $z$ -direction). Hence this radiation can only be observed obliquely (at right angles) to the special direction. It is otherwise in the case of the two transitions  $m \rightarrow m \pm 1$ , which can be observed in all directions. By making observations in the  $x$ -direction, for example, we see the radiation due to the vibration of the dipole in the  $y$ -direction, which in fact is the direction of vibration of the electric vector of the radiation. If we make observations in the  $z$ -direction, we still observe the  $x$ - and  $y$ -components of the radiation, but they are circularly polarized; as we see at once from the discussion above, the transition  $m \rightarrow m + 1$  corresponds to a vibration of the charge-cloud with dipole moment  $x + iy$ , i.e. a circular vibration in the positive direction about the  $z$ -axis; similarly for the transition  $m \rightarrow m - 1$ .

These theoretical statements can be tested directly in the case of the normal Zeeman effect. As is well known, on transverse observation (at right angles to the magnetic field) we see the normal Lorentz triplet, i.e. a splitting-up into three components. Of these the central one, which corresponds to the transition  $m \rightarrow m$  and is therefore not displaced, is polarized in the direction of the magnetic field, while the two other components, corresponding to the transitions  $m \rightarrow m \pm 1$  are transversally polarized. In longitudinal observations the undisplaced component disappears and we see only the two displaced components, which, as theory requires, are circularly polarized.

## XXII. The Formalism of Quantum Mechanics, and the Uncertainty Relation (pp. 84, 135).

The deduction of the uncertainty relation, with the help of diffraction phenomena and other processes capable of being visualized, gives only a result specifying an order of magnitude. To obtain an exact inequality, we must call upon the formalism of quantum mechanics. Of this a short exposition will now be given, and the various theorems will be illustrated by examples in the preceding Appendices.

In wave mechanics, there corresponds to every physical magnitude a (linear) operator  $A$ , which acts upon a wave function  $\psi$  of the co-ordinates  $q$  ( $q$  for brevity representing  $q_1, q_2, \dots$ ), and transforms it into another function  $\phi(q)$ :

$$\phi = A\psi.$$

To the co-ordinate  $q$  itself, considered as an operator, corresponds the operation "multiplication by  $q$ ",  $\phi = q\psi$ ; to the momentum  $p$  corresponds the operator  $\frac{h}{2\pi i} \frac{\partial}{\partial q}$ , so that  $\phi = \frac{h}{2\pi i} \frac{\partial \psi}{\partial q}$ . Generally, we may think of  $A$  as any function of  $q$  and  $p$ ; e.g.  $p^2 + q^2$ , which changes  $\psi$  into the function

$$\phi = (p^2 + q^2)\psi = \left(\frac{h}{2\pi i}\right)^2 \frac{\partial^2 \psi}{\partial q^2} + q^2 \cdot \psi.$$

Two operations  $A$ ,  $B$  are in general not commutative,  $AB \neq BA$ ; e.g.  $pq - qp = h/2\pi i$ . If an operator  $A$  reproduces a function  $\psi$  except for a factor,  $A\psi = a\psi$ , then  $\psi$  is called a proper function of  $A$ , and  $a$  a proper value of  $A$ . An example is given by the equation  $p\psi = a\psi$ , or  $\frac{h}{2\pi i} \frac{\partial \psi}{\partial q} = a\psi$ , with the solution  $\psi = e^{(2\pi i/h)aq}$ , which shows that the proper functions in general are complex numbers. The proper value  $a$  is here any number. The operator  $p$  itself is usually called real, if the proper values  $a$  are taken as real numbers; the operator  $q$  is of course also real. Generally, the following definition is found useful: to every operator  $A$  there corresponds an adjoint  $A^+$  defined by the condition that for any two functions  $\phi$ ,  $\psi$  we are to have

$$\int \phi^* \cdot (A\psi) dq = \int (A^+\phi)^* \cdot \psi dq.$$

If  $A = A^+$ ,  $A$  is said to be real. That is the case, for example, for  $q$  and  $p$ . With  $q$ , this is trivial; with  $p$ , we have, by integration by parts,

$$\begin{aligned} \int \phi^* \cdot (p\psi) dq &= \int \phi^* \left( \frac{h}{2\pi i} \frac{\partial \psi}{\partial q} \right) dq = - \int \psi \frac{h}{2\pi i} \frac{\partial \phi^*}{\partial q} dq \\ &= \int \psi \cdot \left( \frac{h}{2\pi i} \frac{\partial \phi}{\partial q} \right)^* dq = \int \psi \cdot (p\phi)^* dq, \end{aligned}$$

where we assume that the integrated terms vanish. Generally, the proper values of real operators are real; for it follows from  $A = A^+$ , for  $\phi = \psi$ , that the magnitude

$$\int \psi^* (A\psi) dq = \int (A\psi)^* \cdot \psi dq = \left( \int (A\psi) \cdot \psi^* dq \right)^*,$$

and is therefore real. If further  $A\psi = a\psi$ , then, except for the real factor  $\int \psi^* \psi dq$ , the preceding magnitude is equal to  $a$ .

Let  $A$  be a real operator, and  $a_1, a_2$  two of its proper values,  $\psi_1, \psi_2$  the corresponding proper functions, so that we have

$$A\psi_1 = a_1\psi_1, \quad A\psi_2 = a_2\psi_2.$$

Multiply the first equation by  $\psi_2^*$ , the conjugate of the second by  $\psi_1$ ; then integrate over the  $q$ -space and subtract: thus

$$\int [\psi_2^*(A\psi_1) - \psi_1(A\psi_2)^*] dq = (a_1 - a_2) \int \psi_1\psi_2^* dq.$$

The left side vanishes, since  $A = A^+$ . Hence, for any two unequal proper values ( $a_1 \neq a_2$ ),

$$\int \psi_1\psi_2^* dq = 0;$$

the proper functions are said to be orthogonal.

Any proper function can be multiplied by an arbitrary factor; this we determine by the normalizing condition

$$\int \psi_1\psi_1^* dq = \int |\psi_1|^2 dq = 1;$$

thus, for any two proper values  $a_m, a_n$  (assumed to be unequal)

$$\int \psi_m\psi_n^* dq = \delta_{mn},$$

where  $\delta_{mn}$  is equal to 1 or 0, according as  $m = n$ , or  $m \neq n$ . Examples have already been given (oscillator, p. 295; hydrogen atom, p. 298).

A real operator has in general an infinite number of proper values  $a_1, a_2, \dots$ , and proper functions  $\psi_1, \psi_2, \dots$ . The former we here assume for simplicity to be discrete and all different. We can develop any function  $\phi$  in a series of such proper functions (generalized Fourier series):

$$\phi = \sum_n c_n \psi_n.$$

On multiplication by  $\psi_m^*$  this gives, in virtue of the orthogonal and normalizing relations:

$$\int \phi \psi_m^* dq = c_m.$$

Further, we have

$$\begin{aligned} \int |\phi|^2 dq &= \int \phi \cdot \phi^* dq = \sum_n c_n \int \psi_n \phi^* dq \\ &= \sum_n c_n c_n^* = \sum_n |c_n|^2 \end{aligned}$$

We can take for  $\phi$  the function  $B\psi_m$  where  $\psi_m$  is one of the proper

functions of the operator  $A$ , and  $B$  another operator; the coefficients ( $c_n$ ) depend now on both indices  $m$  and  $n$ :

$$B\psi_m = \sum_n B_{nm}\psi_n,$$

$$B_{nm} = \int \psi_n^* (B\psi_m) dq.$$

These  $B_{nm}$  can be considered as the elements of a matrix, representing the operator  $B$  in the system of the proper functions  $\psi_1, \psi_2, \dots$  of the operator  $A$ ; for it can easily be shown that the  $B_{nm}$  have the properties of matrix coefficients. From the definition of the adjoint operator it follows that

$$B_{nm} = \int \psi_n^* (B\psi_m) dq = \int (B^+\psi_n)^* \cdot \psi_m dq;$$

interchanging  $n$  and  $m$  and taking the conjugate we get

$$B_{mn}^* = \int (B^+\psi_m) \cdot \psi_n^* dq = (B^+)_{nm}.$$

This shows that to find the matrix element of the adjoint  $B^+$  one has to interchange lines and columns ( $m \rightleftharpoons n$ ) and to replace each element by the conjugate value.

We consider now the matrix element of the composite operator  $BC$ :

$$(BC)_{nm} = \int \psi_n^* (BC\psi_m) dq = \int \psi_n^* B(C\psi_m) dq;$$

using the definition of the adjoint operator  $B^+$  we get

$$(BC)_{nm} = \int (B^+\psi_n)^* (C\psi_m) dq.$$

Here we substitute the developments

$$B^+\psi_n = \sum_k (B^+)_{kn}\psi_k, \quad C\psi_m = \sum_l C_{lm}\psi_l$$

and find

$$\begin{aligned} (BC)_{nm} &= \int \sum_k (B^+)_{kn}^* \psi_k^* \sum_l C_{lm} \psi_l dq \\ &= \sum_{kl} B_{nk}^* C_{lm} \int \psi_k^* \psi_l dq \\ &= \sum_k B_{nk}^* C_{km}. \end{aligned}$$

This formula shows that the matrix element of the operator product  $BC$  is equal to the element of the matrix product of the matrices belonging to  $B$  and  $C$ . In other words: The operator calculus (or wave mechanics) and the matrix calculus are equivalent representations of quantum theory.

The statistical interpretation of quantum mechanics consists in the following assumptions: To each physical quantity or "observable" belongs a real operator  $A$ . The proper functions  $\psi_1, \psi_2, \dots$  correspond to the quantised states, for which the operator takes the value  $a_1$  or  $a_2$  or  $a_3 \dots$ ; any function  $\phi$  is a state, which is composed of these states, or it may be a distribution of several systems over the states  $a_1, a_2, \dots$ .

The coefficients  $c_n$  of the expansion determine the strength with which the quantum state  $n$  occurs in the general state  $\phi$ . The probability of then finding the proper value  $a_n$  in a measurement is given by

$$w_n = |c_n|^2.$$

If we assume that  $\int |\phi|^2 dq = 1$ , we have

$$\sum_n w_n = \sum_n |c_n|^2 = 1.$$

The mean value of the quantity represented by  $A$  in the state  $\phi$  is:

$$\begin{aligned} \bar{A} &= \int \phi^* \cdot (A\phi) dq = \int \phi^* \sum_n c_n A \psi_n \cdot dq \\ &= \sum_n |c_n|^2 a_n = \sum_n w_n a_n. \end{aligned}$$

The mean value of the product  $AA^+$  is never negative; for, by the definition of  $A^+$ ,

$$\begin{aligned} \overline{AA^+} &= \int \phi^* \cdot (AA^+\phi) dq = \int \phi^* \cdot A(A^+\phi) dq \\ &= \int (A^+\phi)^* \cdot (A^+\phi) dq = \int |A^+\phi|^2 dq \geq 0. \end{aligned}$$

We can now deduce inequalities referring to the mean values of two real operators  $A, B$ —inequalities which lead to Heisenberg's uncertainty relation.

From the definition of  $A^+$ , it follows on multiplication by  $i$  that

$$\int \phi^* \cdot (iA\psi) dq = - \int (iA^+\phi)^* \cdot \psi dq,$$

that is to say,  $(iA)^+ = -iA^+$ . Somewhat more generally, we have also

$$(A + iB)^+ = A^+ - iB^+.$$

If  $A, B$  are real, and  $\lambda$  is a real number, then

$$\overline{(A + i\lambda B)(A + i\lambda B)^+} \geq 0,$$

or

$$\overline{(A + i\lambda B)(A - i\lambda B)} = \bar{A}^2 + \bar{B}^2 \lambda^2 - i(\overline{AB - BA}) \lambda \geq 0.$$

From this it follows that  $\overline{AB - BA}$  is purely imaginary. The minimum of the last expression occurs when

$$\lambda = \frac{i}{2} \frac{\overline{AB - BA}}{B^2},$$

and is equal to

$$\overline{A^2} + \frac{1}{4} \frac{(\overline{AB - BA})^2}{B^2} \geq 0.$$

Hence

$$\overline{A^2} \cdot \overline{B^2} \geq -\frac{1}{4} (\overline{AB - BA})^2.$$

Now replace  $A$  and  $B$  by  $\delta A = A - \bar{A}$  and  $\delta B = B - \bar{B}$ ; then

$$\begin{aligned} \delta A \delta B - \delta B \delta A &= AB - A\bar{B} - \bar{A}B + \bar{A}\bar{B} \\ &\quad - BA + \bar{B}A + B\bar{A} - \bar{B}\bar{A} \\ &= AB - BA, \end{aligned}$$

and the preceding equation gives

$$(\overline{\delta A})^2 \cdot (\overline{\delta B})^2 \geq -\frac{1}{4} (\overline{AB - BA})^2.$$

If  $p$  and  $q$  are conjugate operators (momentum and co-ordinate), we have

$$pq - qp = \frac{h}{2\pi i};$$

therefore

$$(\overline{\delta p})^2 \cdot (\overline{\delta q})^2 \geq \frac{h^2}{16\pi^2}.$$

For the root-mean-square deviations

$$\Delta p = \sqrt{(\overline{\delta p})^2}, \quad \Delta q = \sqrt{(\overline{\delta q})^2}$$

we therefore have

$$\Delta p \Delta q \geq \frac{h}{4\pi}.$$

The sign of equality holds only for one definite distribution (proper function), viz. the Gaussian error function, which occurs in the case of the linear oscillator (see Appendix XII, p. 282; Appendix XVI, p. 295; and Appendix XXXII, p. 343).

**XXIII. Anomalous Zeeman Effect in the D Lines of Sodium (p. 146).**

We shall now deduce the splitting pattern of the  $D$  lines of sodium in the anomalous Zeeman effect. As we stated at p. 139, the  $D_1$  line corresponds to the transition from a  $p$  term with inner quantum number  $\frac{1}{2}$ , i.e. from ( $l = 1, j = \frac{1}{2}$ ), to an  $s$  term, i.e. to ( $l = 0, j = \frac{1}{2}$ ); the  $D_2$  line represents a transition from ( $l = 1, j = \frac{3}{2}$ ) to ( $l = 0, j = \frac{1}{2}$ ).

We begin by determining Landé's splitting factors for the three terms in question. We have  $s = \frac{1}{2}$ ; the formula

$$g = \frac{3}{2} + \frac{s(s+1) - l(l+1)}{2j(j+1)}$$

then gives the following results:

$$l = 0, j = \frac{1}{2}: g = \frac{3}{2} + \frac{\frac{1}{2} \times \frac{3}{2}}{2 \times \frac{1}{2} \times \frac{3}{2}} = 2;$$

$$l = 1, j = \frac{1}{2}: g = \frac{3}{2} + \frac{\frac{1}{2} \times \frac{3}{2} - 1 \times 2}{2 \times \frac{1}{2} \times \frac{3}{2}} = \frac{2}{3};$$

$$l = 1, j = \frac{3}{2}: g = \frac{3}{2} + \frac{\frac{1}{2} \times \frac{3}{2} - 1 \times 2}{2 \times \frac{3}{2} \times \frac{5}{2}} = \frac{4}{3}.$$

In the two following diagrams we collect the values of the separations of the terms, taking the separation for the normal Zeeman effect as unit. That is, we write down the values  $mg$  for the upper and lower terms of the two lines. The values of  $m$ , like  $j$ , must be halves of odd numbers, as they are equal to  $-j, -j+1, \dots, j$ .

$D_1$	$m =$	$-\frac{1}{2}$	$+\frac{1}{2}$
$l = 1 \quad j = \frac{1}{2}$		$-\frac{1}{3}$	$+\frac{1}{3}$
$\downarrow$			
$l = 0 \quad j = \frac{1}{2}$		$-1$	$+1$

$D_2$	$m =$	$-\frac{3}{2}$	$-\frac{1}{2}$	$\frac{1}{2}$	$\frac{3}{2}$
$l = 1 \quad j = \frac{3}{2}$		$-2$	$-\frac{2}{3}$	$+\frac{2}{3}$	$2$
$\downarrow$					
$l = 0 \quad j = \frac{1}{2}$			$-1$	$+1$	

The arrows indicate the possible transitions, i.e. give the positions of the lines in the Zeeman effect. Here the selection rules for the magnetic quantum number  $m$  must be taken into account. These can be deduced from the correspondence principle in exactly the same way as at p. 110 (see also Appendix XXI, (p. 308)). As  $m$  denotes the precessional motion about the direction of the field, the transition  $\Delta m = \pm 1$  corresponds to the classical vibrations at right angles to  $H$ ;



these components of the radiation are called  $\sigma$  components. When observed in the longitudinal direction (in the direction of the field), they appear circularly polarized (as predicted by the classical theory); when observed in the transverse direction, they appear linearly polarized at right angles to the field. Further, the transitions  $\Delta m = 0$  are also permissible; they correspond to the classical vibrations in the direction of the field. These, known as  $\pi$  components, are visible as vibrations parallel to the field, on transverse observation only (in the direction of vibration of a dipole, i.e. in the direction of the magnetic field in the present case, the radiation is zero).

These selection rules at once determine the positions of the components in the Zeeman splitting of the  $D$  lines. We measure their displacements on either side from a central zero position, as before taking the separation in the normal Zeeman effect, i.e.  $\nu_L$  in the frequency scale, as unit. The  $\pi$  components are shown above the horizontal axis and the  $\sigma$  components below. We thus obtain the splitting diagram of fig. 6,

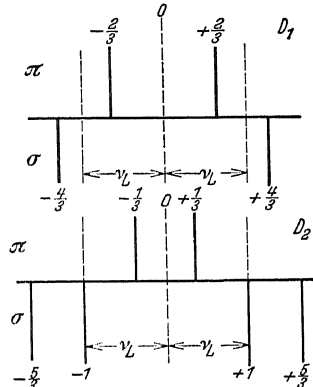


Fig. 6.—Splitting of the sodium  $D$  lines in the anomalous Zeeman effect. The splitting ( $\nu_L$ ) in the normal effect is shown as unity. The ( $\pi$ -) components polarized parallel to the field are shown above; the ( $\sigma$ -) components polarized perpendicular to the field, below.

which is found to be in complete agreement with that obtained experimentally (see fig. 2b, Plate VII).

#### XXIV. Enumeration of the Terms in the case of two $p$ -Electrons (p. 160).

Here we consider an example of the enumeration of the terms of an atom in the case of two valency electrons (Hund); a knowledge of the number of terms is of great importance in connexion with the analysis of the corresponding spectrum and is also required in problems of a statistical nature.

For example, we suppose that the values of the terms are determined by two  $p$ -electrons and inquire what the number of terms is. Here we must consider whether these two electrons are equivalent or not, i.e. whether their principal quantum numbers are the same or not; we have of course assumed that their azimuthal quantum numbers are both equal to 1.

We shall take the second case first as it is the simpler; here we do not need to trouble about the exclusion principle, seeing that the principal quantum numbers are different. According to the rules for combining angular momenta in the quantum theory, the resultant orbital momentum  $l$  may have the three values 0, 1, 2 ( $l_1$  and  $l_2$  may be parallel ( $l=2$ ) or anti-parallel ( $l=0$ ) or may be inclined to one another at an angle such that their vector sum is 1). Combination of the two spin moments gives resultant spin 1 for the parallel position, resultant spin 0 for the anti-parallel position. Hence both a triplet system and a singlet system occur. Moreover, owing to the three possibilities for  $l$  there are  $S$ ,  $P$ , and  $D$  terms. The following terms accordingly appear:

$$^1S, ^1P, ^1D; ^3S, ^3P, ^3D.$$

As the  $^3S$  term is single (see the discussion in the text (pp. 139 and 158)), in this case there are ten different terms in all.

If an external field is applied, however, this number is considerably increased owing to the various settings of the resultant angular momentum relative to the special direction; a term with resultant angular momentum  $j$  has  $2j+1$  possible settings in the field. We thus obtain the following scheme:

Term	$l$	$s$	$j$	$2j+1$
$^1S$	0	0	0	1
$^1P_1$	1	0	1	3
$^1D_2$	2	0	2	5
$^3S$	0	1	1	3
$^3P_0$	1	1	0	1
$^3P_1$	1	1	1	3
$^3P_2$	1	1	2	5
$^3D_1$	2	1	1	3
$^3D_2$	2	1	2	5
$^3D_3$	2	1	3	7
Total ..				36

In an external field there are accordingly 36 different energy levels in all, for the case of two non-equivalent  $p$ -electrons.

The enumeration becomes more complicated in the case where the electrons are equivalent, as here the exclusion principle must be taken into account. Hence it is necessary to write down the complete system of eight quantum numbers and strike out all the cases where all the quantum numbers of the two electrons are the same. In the first instance, therefore, we imagine the directional degeneracy removed

by an external field, so that there is a meaning in stating the components of the angular momenta in the special direction. As quantum numbers we shall in addition to  $n_1 = n_2$  and  $l_1 = l_2 = 1$  use  $\mu_1, \mu_2$ , the projections of the orbital momenta, and  $\sigma_1, \sigma_2$ , the projections of the spin momenta, along the special direction.

In the following table we have collected all the possible combinations for these four quantum numbers, omitting all those combinations which would contradict the exclusion principle. Further, we have not written down those combinations which arise from a given set merely by interchanging the quantum numbers of the two electrons; for since we cannot distinguish one electron from the other these of course are identical and are associated with the same energy term. Hence we have the following table:

$\mu_1$	$\mu_2$	$\sigma_1$	$\sigma_2$	$\mu$	$\sigma$
1	1	$\frac{1}{2}$	$-\frac{1}{2}$	2	0
1	0	$\pm\frac{1}{2}$	$\pm\frac{1}{2}$	1	1, 0, 0, -1
1	-1	$\pm\frac{1}{2}$	$\pm\frac{1}{2}$	0	1, 0, 0, -1
0	0	$\frac{1}{2}$	$-\frac{1}{2}$	0	0
0	-1	$\pm\frac{1}{2}$	$\pm\frac{1}{2}$	-1	1, 0, 0, -1
-1	-1	$\frac{1}{2}$	$-\frac{1}{2}$	-2	0

In the fifth and sixth columns we have inserted the sums  $\mu = \mu_1 + \mu_2$  and  $\sigma = \sigma_1 + \sigma_2$ . We see, therefore, that in the magnetic field there are only 15 terms in all (magnetic splitting), while, as we saw above, in the case of non-equivalent electrons there are 36 terms in all in the magnetic field.

We are, however, not so much interested in the splitting of the terms in the magnetic field as in values for the undisturbed atom. We must therefore combine into one term of the undisturbed atom those terms in the magnetic field which have the same inner quantum number  $j$  and the same orbital angular momentum  $l$ ; for we know from the above that one term with inner quantum number  $j$  is split up in a magnetic field into  $2j + 1$  terms. Now from the above table we see that there must be at least one term with orbital angular momentum  $l = 2$  (a  $D$  term), as there are components of this momentum with values 2 and -2 in the specified direction. We see at the same time that the corresponding spin quantum number  $s$  must vanish, as in these terms the spin component  $\sigma = 0$  alone appears; we accordingly have a  $^1D$  term with inner quantum number  $j = \vec{l} + \vec{s} = 2$ , which

must split up into five terms in the magnetic field, terms for which  $\sigma = 0$  and  $\mu = -2, -1, \dots, +2$ . Of the remaining ten terms, we may likewise readily convince ourselves that nine correspond to  $P$  terms ( $l = 1$ ) with spin quantum numbers  $\sigma = -1, 0, 1$ , while the last term is an  $^1S$  term. We accordingly have the following arrangement:

Multiplet		Term Notation	
$\mu = -2, -1, 0, 1, 2$	corresponds to	$l = 2$	$^1D$
$\sigma = 0$	„	$s = 0$	
$\mu = -1, 0, 1$	„	$l = 1$	$^3P$
$\sigma = -1, 0, 1$	„	$s = 1$	
$\mu = 0$	„	$l = 0$	$^1S$
$\sigma = 0$	„	$s = 0$	

Here (in the case of equivalent  $p$ -electrons), therefore, there are only five terms in the absence of a magnetic field, as compared with ten terms in the case of non-equivalent electrons. The terms  $^3S$ ,  $^1P$ , and  $^3D$  found above fall out here (owing to the exclusion principle).

An analogous enumeration for the case of a greater number of equivalent  $p$ -electrons is given in the following table. As regards the

Number of $p$ Electrons	Terms
6	$^1S$
1 or 5	$^3P$
2 or 4	$^1S, ^3P, ^1D$
3	$^4S, ^2P, ^2D$

energies of the terms, a quantum-mechanical estimation agrees with experiment in giving the result that the term with the greatest multiplicity, i.e. with the greatest value of  $s$ , is always the lowest; that is, if there are two or four  $p$ -electrons the ground state is a  $^3P$  term, while if there are three  $p$ -electrons it is a  $^4S$  term. If there are several terms of the same multiplicity, that with the greatest value of  $l$  is the lowest.

## XXV. Temperature Variation of Paramagnetism (p. 166).

In order to understand the variation of paramagnetism with temperature we shall consider a simplified model of a paramagnetic substance. We think of the substance as consisting of a large number of particles, all of which have the same magnetic moment  $M$ . We shall also in the first instance ignore azimuthal quantisation, i.e. we assume that this moment can be inclined at *any* angle  $\theta$  to the direction of the field, and not merely at certain specified angles.

So long as a particle is moving undisturbed in the magnetic field, the magnetic moment, which is always associated with an angular momentum about the same axis (cf. the top), will execute a precessional motion about the direction of the field, the angle at which it is inclined to the field remaining constant. The magnetic energy for this orientation is given by

$$E = -MH \cos \theta.$$

As a result of interaction with the other particles (collisions), however, this state of equilibrium is disturbed, the angle of inclination after collision being different from that before collision. From the kinetic theory we already know that the collisions due to thermal motions have the effect that in course of time the particles become uniformly distributed among all possible states (here possible angles of inclination to the field). The magnetic field, however, acts against this equalization; accordingly, settings of the moment in its direction are favoured as regards energy compared with those in the opposite direction. A state of equilibrium will be established, which can be determined by statistical methods.

In our discussion of the kinetic theory of gases we have already shown that a statistical argument gives Boltzmann's distribution law for the probability of a definite energy state, according to which a definite state with energy  $E$  has the probability

$$W \sim e^{-E/kT}.$$

The form in which the temperature  $T$  appears in this expression is deduced from thermodynamical considerations;  $k$  is Boltzmann's constant. In our case, therefore, the probability that the magnetic moment will set itself at a definite angle  $\theta$  to the direction of the field is given by

$$W \sim e^{MH \cos \theta / kT} = e^{\beta \cos \theta} \quad (\beta = MH/kT).$$

We now investigate the mean magnetic moment in the direction of the field. If the field were equal to zero or if the thermal motion (in the case of very high temperatures) greatly preponderated over the directional effect of the magnetic field, so that the distribution of directions became almost uniform, the mean magnetic moment in a specified direction, in the direction of the field in particular, would be approximately or accurately equal to zero. At low temperatures, or with magnetic fields so strong that the magnetic energy  $MH$  is of the same order as the thermal energy  $kT$ , there will be a marked preference for

the direction of the field, giving rise to a finite mean magnetic moment in this direction. If we adopt the classical ideas, according to which all directions of the moment are equally possible, the calculation is easily carried out. By definition

$$\overline{M \cos \theta} = M \frac{\int_0^\pi \cos \theta e^{\beta \cos \theta} \sin \theta d\theta}{\int_0^\pi e^{\beta \cos \theta} \sin \theta d\theta}.$$

Evaluating the integrals, we obtain

$$\begin{aligned} \overline{M \cos \theta} &= M \frac{d}{d\beta} \log \int_0^\pi e^{\beta \cos \theta} \sin \theta d\theta = M \frac{d}{d\beta} \log \frac{e^\beta - e^{-\beta}}{\beta} \\ &= M \left( \frac{e^\beta + e^{-\beta}}{e^\beta - e^{-\beta}} - \frac{1}{\beta} \right) \\ &= M \left( \coth \beta - \frac{1}{\beta} \right). \end{aligned}$$

In the limiting case  $\beta \ll 1$ , i.e. for weak fields or high temperatures, we obtain the expression

$$\overline{M \cos \theta} = M \left( \frac{1}{3} \beta + \dots \right) = \frac{M^2 H}{3kT}$$

by expanding in powers of  $\beta$ .

In this limiting case, however, we could have obtained the result more easily by introducing a power series in  $\beta$  into the formula defining  $\overline{M \cos \theta}$ . We then have

$$\overline{M \cos \theta} = M \frac{\int_0^\pi \cos \theta (1 + \beta \cos \theta + \dots) \sin \theta d\theta}{\int_0^\pi (1 + \beta \cos \theta + \dots) \sin \theta d\theta}.$$

The first term in the numerator vanishes on integration, while the second gives  $2\beta/3$ . Even the first term in the denominator differs from zero, its value being 2. The quotient gives the expression stated above.

Now it is important to know that the same formula results from calculations on the basis of the quantum theory, i.e. if we take into account only a finite number of positions of the moment. We assume that  $\beta$  has small values and the resultant angular momentum large

values. The fact that in this case we obtain the classical result for  $\overline{M \cos \theta}$  is understandable if we consider the correspondence principle (the limiting case of large quantum numbers). It is perhaps advisable, however, for the reader to convince himself that this statement is correct by direct evaluation of the sums. If the resultant angular momentum is  $j$ , there are  $2j + 1$  possible settings of the moment relative to the direction of the field; in fact, the projection of  $j$  on this can have the values  $-j, -j + 1, \dots, +j$ . In the previous calculation we therefore have merely to replace  $\cos \theta$  by  $m/j$  and replace the integrals by sums:

$$\overline{M \cos \theta} = M \frac{\sum_{-j}^{+j} \frac{m}{j} e^{\beta m/j}}{\sum_{-j}^{+j} e^{\beta m/j}}.$$

Expanding in powers of  $\beta$  and making use of the formula

$$\sum_0^n \nu^2 = \frac{n(n+1)(2n+1)}{6},$$

we have

$$\begin{aligned} \overline{M \cos \theta} &= M \frac{\sum_{-j}^{+j} \frac{m}{j} \left(1 + \frac{\beta m}{j}\right)}{\sum_{-j}^{+j} \left(1 + \frac{\beta m}{j}\right)} = M \beta \frac{\frac{1}{3} \frac{(j+1)(2j+1)}{j}}{2j+1} \\ &= \frac{M \beta}{3} \frac{j+1}{j}, \end{aligned}$$

or, for large values of  $j$ ,  $\overline{M \cos \theta} = \frac{M \beta}{3}$ , as we previously obtained by classical methods.

The susceptibility  $\chi$  is defined as the magnetic moment of the substance in question per unit field intensity  $H$ ; per mole

$$\chi = \frac{LM^2}{3kT} = \frac{(LM)^2}{3RT}.$$

This is known as Curie's law; the magnetic susceptibility is inversely proportional to the temperature, i.e. decreases with increasing temperature.

We conclude with a few brief remarks on the simplifications made above. The account which we have just given certainly suffices for a rough survey of the variation of paramagnetism with temperature;

but if, in particular, we wish to draw conclusions about the magnitude of the atomic magnetic moments from the measured values of the susceptibility, more detailed arguments become necessary. In the first place, as we have already indicated, the existence of azimuthal quantisation must be taken into account; this has the effect of introducing the factor  $(j + 1)/j$  into the expression given above for the mean magnetic moment in the direction of the field. Further, we must take account of the fact that in general all the atoms of a substance do not have the same magnetic moment. When discussing the anomalous Zeeman effect we saw that the resultant magnetic moment of an atom is equal to the Bohr magneton multiplied by the resultant angular momentum  $j$  and the so-called Landé factor  $g$ , which depends on the three quantum numbers  $j$ ,  $s$ , and  $l$ . One would therefore have to take the average for all possible combinations of the quantum numbers.

In this connexion we add the following remark. Weiss imagined that his measurements of the susceptibilities of various substances necessitated the conclusion that an elementary magnetic moment or "magneton" does exist, of which the magnetic moments of the various substances are integral multiples; the Weiss magneton is about one-fifth of the Bohr magneton  $\frac{eh}{4\pi\mu c}$  whose existence follows from the theory. This was due to his application to the experimental results of a formula which had been deduced without reference to the quantum theory. Discussion of more recent results in the light of the quantum theory has established the Bohr magneton as an elementary unit.

#### XXVI. Theory of Nuclear Disintegration (p. 182).

Let the potential for an  $\alpha$ -particle, which has been emitted from a nucleus of atomic number  $Z$  and is therefore in the field of the residual nucleus  $Z - 2$ , be  $V(r)$ . For great distances this is the Coulomb potential, i.e.

$$V(r) = \frac{2(Z - 2)e^2}{r}, \text{ for } r > r_0;$$

for  $r < r_0$  the form  $V(r)$  is unknown, but special assumptions about it are unnecessary, apart from this, that it must have the crater-like character represented in fig. 18, p. 182.

According to Laue, the frequency  $\lambda$  of the emission can be split up into two factors. We think of the  $\alpha$ -particle as oscillating to and fro in the crater, so that it strikes the wall  $n$  times per second. At



each collision there is a certain probability  $p$  that it passes the wall. Hence

$$\lambda = np.$$

The order of magnitude of  $n$  we can take to be  $v/2r_0$ , where  $v$  is the velocity of the  $\alpha$ -particle; the latter in its turn can be determined by putting the wave-length  $h/mv$  of the associated de Broglie wave equal to  $2r_0$ . We thus find

$$n = \frac{h}{4mr_0^2}.$$

To calculate  $p$ , we have to find a suitable solution of Schrödinger's equation

$$\frac{h^2}{8\pi^2m} \frac{d^2\psi}{dr^2} + (E - V(r))\psi = 0.$$

For  $r < r_0$ , the function  $\psi$  will oscillate in some way or other; for large values of  $r$  it will be a progressive wave. If the energy of the  $\alpha$ -particle is  $E$ , the latter state will be attained at about the distance where

$$E = \frac{2(Z-2)e^2}{r_1}, \quad r_1 = \frac{2(Z-2)e^2}{E}.$$

In the intermediate zone from  $r_0$  to  $r_1$ ,  $\psi$  will fall off exponentially. Thus, clearly we have approximately

$$p = \left| \frac{\psi(r_1)}{\psi(r_0)} \right|^2.$$

But Coulomb's law still holds in this intermediate zone; we have therefore to solve the equation

$$\frac{h^2}{8\pi^2m} \psi'' + \left( E - \frac{2(Z-2)e^2}{r} \right) \psi = 0.$$

Since  $Z$  is a large number, the following approximate method gives what is wanted. Put  $\psi = e^{(2\pi/h)y(r)}$ ; we then obtain the equation

$$\frac{h}{2\pi} y'' + y'^2 - F(r) = 0,$$

where

$$F(r) = 2m \left( -E + \frac{2(Z-2)e^2}{r} \right).$$

Neglecting the term multiplied by  $\hbar$ , we get

$$y' = \sqrt{F(r)}, \quad y = \int_a^r \sqrt{F(r)} dr,$$

and therefore

$$\begin{aligned} \frac{\psi_1}{\psi_0} &= e^{(2\pi/\hbar)(y(r_1) - y(r_0))} \\ &= e^{(2\pi/\hbar) \int_{r_0}^{r_1} \sqrt{F(r)} dr}. \end{aligned}$$

If we substitute the expression for  $F(r)$ , the integration can be carried out, and we find

$$p = e^{-(2n_0 - \sin 2n_0)(8\pi e^2/\hbar)(Z-2)/v},$$

where  $n_0$  is given by the equation

$$\cos^2 n_0 = \frac{r_0 E}{2(Z-2)e^2}.$$

If we expand the exponent in powers of the quantity on the right, we obtain finally the approximation

$$p = e^{-8\pi^2 e^2 (Z-2)/(\hbar v) + (16\pi e \sqrt{m/\hbar})(\sqrt{Z-2})r_0}.$$

The disintegration constant  $\lambda$  is therefore given by

$$\log \lambda = \log \frac{\hbar}{4\pi m r_0^2} - \frac{8\pi^2 e^2 (Z-2)}{\hbar v} + \frac{16\pi e}{\hbar} \sqrt{m(Z-2)} r_0,$$

or, numerically,

$$\log_{10} \lambda = 20.46 - 1.191 \times 10^9 \frac{Z-2}{v} + 4.084 \times 10^6 \sqrt{(Z-2)} r_0;$$

here  $\lambda$  is measured in  $\text{sec.}^{-1}$ , and is connected with the half-value period by the equation

$$T = \frac{\log_e 2}{\lambda} = \frac{0.6931}{\lambda}.$$

The law thus obtained differs from the empirical formula of Geiger and Nuttall in this respect, that it is not linear in  $v$ , but in  $1/v$ ; however, since the variation of  $v$  is confined to comparatively narrow limits, from  $v = 1.4 \times 10^9$  cm./sec. to  $v = 2.0 \times 10^9$  cm./sec., the difference is slight. In consequence of the large factor in the second term—in which  $v$  appears—the range of values of the disintegration constant is extremely wide.

# XXVII. The Stefan-Boltzmann Law and Wien's Displacement Law (p. 187).

The thermodynamical proof of the Stefan-Boltzmann law rests on the existence of radiation pressure. We imagine an enclosure shut off by a movable piston with a reflecting surface. The radiation field exerts a pressure on the piston; its magnitude is a function of  $u$ , the energy density of the radiation in the enclosure. In fact, both Maxwell's theory and the quantum (corpuscular) theory of light give the formula

$$p = \frac{1}{3}u.$$

This radiation pressure is due to the momentum which the radiation carries with it. In the case of the quantum theory this is clear, for according to this theory each light quantum of energy  $h\nu$  possesses momentum  $h\nu/c$ . Maxwell's theory also ascribes to every radiation field, with energy-density  $u$ , the "momentum density"  $|\mathbf{g}| = \frac{u}{c} = \frac{1}{c^2}|\mathbf{S}|$ ,

where  $\mathbf{S}$  is the Poynting vector; this is proved, e.g., by imagining a plane light wave to fall on a metal, where it is absorbed, and calculating the resulting force on the metal from Maxwell's equations. On both theories, then, the momentum contained in a definite volume of the radiation field is equal to the radiation energy contained in this volume divided by  $c$ . The rest of the calculation of the radiation pressure is now just like that of the mechanical pressure in the kinetic theory of gases. In the interval of time  $dt$  radiation energy  $\frac{uc}{4\pi} dt d\omega \cos \theta$

falls on one square centimetre of the boundary from the solid angle  $d\omega$ ; on reflection at the boundary it transfers to the latter a momentum equal to twice the component of the radiation momentum normal to the wall, i.e. the momentum  $2 \frac{u}{c} \frac{cdt}{4\pi} d\omega \cos^2 \theta$ . If we integrate this expression with respect to  $\theta$  and  $\phi$ , where  $d\omega = \sin \theta d\theta d\phi$ , and  $\theta$  goes from 0 to  $\frac{1}{2}\pi$ ,  $\phi$  from 0 to  $2\pi$ , we obtain the pressure  $p = \frac{1}{3}u$  as the momentum transferred to the boundary per unit time as a result of the reflection of the radiation.

We now regard the radiation as a thermodynamic engine, and apply the fundamental equation of thermodynamics, which includes both the first law and the second. If  $W$  is the total energy,  $S$  the entropy,  $T$  the absolute temperature, and  $V$  the volume, then, as we know,

$$TdS = dW + pdV.$$

Now  $W = Vu$ , where the energy-density  $u$  is a function of  $T$  alone; hence

$$TdS = u dV + V \frac{du}{dT} dT + \frac{1}{3} u dV = V \frac{du}{dT} dT + \frac{4}{3} u dV.$$

From this we may conclude that

$$\frac{\partial S}{\partial V} = \frac{4}{3} \frac{u}{T}, \quad \frac{\partial S}{\partial T} = \frac{V}{T} \frac{du}{dT},$$

and hence

$$\frac{\partial}{\partial V} \left( \frac{V}{T} \frac{du}{dT} \right) = \frac{\partial^2 S}{\partial T \partial V} = \frac{4}{3} \frac{d(u/T)}{dT},$$

or

$$\frac{1}{T} \frac{du}{dT} = \frac{4}{3} \left( \frac{1}{T} \frac{du}{dT} - \frac{u}{T^2} \right).$$

Accordingly

$$\frac{du}{dT} = 4 \frac{u}{T} \quad \text{or} \quad u = aT^4,$$

which is the Stefan-Boltzmann radiation law. It is also easy to verify that the entropy  $S$  is given by  $S = \frac{4}{3} aVT^3$ .

Wien's displacement law depends on the existence of the Doppler effect. As is well known, to an observer at rest a wave motion whose source is moving appears to have its frequency altered. In actual fact, it is only the component of the velocity in the line of observation that is effective, and we have the formula  $\frac{\Delta\nu}{\nu} = \frac{v}{c} \cos \theta$ , where  $\nu$  is the frequency,  $\Delta\nu$  the change of frequency,  $v$  the velocity of the source, and  $\theta$  the angle between the direction in which the source is moving and the line of observation.

Now if a mirror at which a wave of frequency  $\nu$  is reflected moves with velocity  $v$  in the direction in which the light is propagated, we may think of the incident wave as coming from a source of light at rest. Then the reflected wave must behave as if it came from the mirror image of this imaginary source of light. Owing to the motion of the mirror, however, this mirror image moves in the direction of the normal to the mirror with velocity  $2v$ . Hence the change of frequency due to reflection at the moving mirror is  $-2 \frac{v}{c} \nu$ , and for oblique incidence at the angle  $\theta$  we have

$$\nu' = \nu \left( 1 - \frac{2v}{c} \cos \theta \right).$$

Further, we can easily see that the intensity of radiation due to a train of waves falling on the moving mirror is altered by reflection in the same ratio. For the energy of radiation falling on the whole mirror of area  $A$  in time  $dt$  is  $IA dt$ , and that reflected is  $I'A dt$ . The difference must be equal to the work done by radiation pressure  $P$  in moving the mirror, i.e.  $PAvdt$ ; but, as shown above,  $P = \frac{2I}{c} \cos \theta$ . We thus obtain

$$I' = I \left( 1 - \frac{2v}{c} \cos \theta \right).$$

We shall now apply thermodynamics not to the total radiation but to a definite narrow range of wave-length. Here we must take note of the fact that whenever work is done in connexion with a motion of the reflecting piston the radiation will be displaced into another region of frequency according to the Doppler formula. This displacement does not by any means vanish when we make the mirror move infinitely slowly. To convince ourselves of this, we imagine the enclosure entirely surrounded by reflecting walls, so that a particular beam of light will go on moving in a zigzag path in the enclosure, meeting the mirrors again and again. If the velocity of the piston is now reduced to half, the number of reflections at it per unit time will be just twice as great, so that the total Doppler displacement as  $v \rightarrow 0$  tends to a finite limit.

We shall now calculate  $\int u_\nu d\nu$ , the total change of energy of the radiation in an enclosure corresponding to a definite range of frequency ( $\nu, \nu + d\nu$ ) during time  $dt$ , arising from reflection at the mirror moving with velocity  $v$ . In the first place, all components of the radiation of this region of the spectrum which fall on the mirror are removed from the region in question owing to the Doppler effect. On the other hand, all components of the radiation which before reflection lay in the interval between  $\nu'$  and  $\nu' + d\nu'$  will now be brought into the region if  $\nu' = \nu(1 + 2\frac{v}{c} \cos \theta)$ . The quantity of energy from the range of frequency ( $\nu, \nu + d\nu$ ) falling from solid angle  $d\omega$  on  $A$ , the area of the mirror, per second is

$$\Delta E = u_\nu \frac{c}{4\pi} A \cos \theta dt d\omega d\nu.$$

From this we obtain the energy thrown into the frequency range ( $\nu, d\nu$ ) on reflection by (1) multiplying by the energy factor  $(1 - \frac{2v}{c} \cos \theta)$ ,

(2) replacing  $d\nu$  by  $d\nu' = d\nu(1 + \frac{2v}{c} \cos \theta)$ , and (3) introducing the expansion

$$u_{\nu'} = u_{\nu(1+2v/c \cos \theta)} = u_{\nu} + \nu \frac{\partial u_{\nu}}{\partial \nu} 2 \frac{v}{c} \cos \theta.$$

The product of the first two factors, however, differs from  $d\nu$  merely by a term of the second order in  $v/c$ , which can be omitted. For the energy diverted into the range  $(\nu, d\nu)$  by reflection we accordingly have

$$\left( u_{\nu} + \nu \frac{\partial u_{\nu}}{\partial \nu} 2 \frac{v}{c} \cos \theta \right) \frac{c}{4\pi} A \cos \theta dt d\omega d\nu;$$

the increase in the energy of the beam as a result of reflection is therefore

$$\nu \frac{\partial u_{\nu}}{\partial \nu} \frac{1}{2\pi} dV d\nu \cos^2 \theta d\omega,$$

since  $A\nu dt = dV$ . Integrating over the hemisphere (for which  $\int \cos^2 \theta d\omega = 2\pi/3$ ), we obtain the total increase of energy due to reflection,  $d(u_{\nu} V) d\nu$ . Hence, for constant  $\nu$ ,

$$d(u_{\nu} V) = \frac{1}{3} \nu \frac{\partial u_{\nu}}{\partial \nu} dV.$$

This is a differential equation for  $u_{\nu}$  as a function of  $\nu$  and  $V$ :

$$\frac{\partial(Vu_{\nu})}{\partial V} = \frac{1}{3} \frac{\partial u_{\nu}}{\partial \nu} \nu \quad \text{or} \quad V \frac{\partial u_{\nu}}{\partial V} = \frac{1}{3} \nu \frac{\partial u_{\nu}}{\partial \nu} - u_{\nu}.$$

We easily see that this equation is satisfied by the relation

$$u_{\nu} = \nu^3 \phi(\nu^3 V),$$

where  $\phi$  is an arbitrary function.

We now take a step further. We imagine that the change of volume is adiabatic, i.e. takes place without heat being supplied. This means that the entropy of the radiation in the enclosure remains constant during the compression. Now we saw above in deducing the Stefan-Boltzmann law that this entropy is proportional to the product of the volume  $V$  and the cube of the temperature, so that the constancy of the entropy implies that

$$VT^3 = \text{const.}$$

If in order to make the law of radiation independent of the size and

shape of the enclosure we introduce the temperature instead of  $V$  by means of this relation, we have

$$u_\nu = \nu^3 F\left(\frac{\nu}{T}\right),$$

which is Wien's displacement law, given in the text.

### XXVIII. Absorption by an Oscillator (p. 189).

We shall now add the proof of the formula

$$\delta W = \frac{1}{3} \frac{\pi e^2}{m} u_\nu$$

given in the text for the work done per second by a radiation field on an oscillator. We define the radiation field by stating the relationship between the electric vector  $E$  and the time. For reasons of convergence we shall assume that the radiation field exists only between the instants  $t = 0$  and  $t = T$ ; it is easy to pass to the limiting case  $T \rightarrow \infty$  subsequently. We now express one component, e.g.  $E_x$ , by its spectrum (Fourier integral),

$$E_x(t) = \int_{-\infty}^{+\infty} f(\nu) e^{2\pi i \nu t} d\nu,$$

where the amplitudes  $f(\nu)$  are determined by

$$f(\nu) = \int_0^T E_x(t) e^{-2\pi i \nu t} dt,$$

$E_x(t)$  being zero outside the interval  $t = 0$  to  $t = T$ ; as  $E_x$  is real, the conjugate complex quantity satisfies the equation

$$f^*(\nu) = f(-\nu).$$

Analogous relationships hold for the two other components. According to the laws of electrodynamics, the total energy-density of the radiation field is then given by

$$u = \frac{1}{8\pi} (\overline{E^2} + \overline{H^2}) = \frac{1}{4\pi} \overline{E^2} = \frac{3}{4\pi} \overline{E_x^2},$$

the last of these relations follows from symmetry. The bars denote time-averages. Further,

$$\overline{E_x^2} = \frac{1}{T} \int_0^T E_x^2 dt = \frac{1}{T} \int_0^T E_x dt \int_{-\infty}^{+\infty} f(\nu) e^{2\pi i \nu t} d\nu,$$

if we replace one factor by the above Fourier expression. If we now change the order of integration, we immediately obtain, as the value of the new integral with respect to  $t$  has been shown to be  $f^*(\nu)$ ,

$$\begin{aligned}\overline{E_x^2} &= \frac{1}{T} \int_{-\infty}^{+\infty} f(\nu) d\nu \int_0^T E_x e^{2\pi i \nu t} dt \\ &= \frac{1}{T} \int_{-\infty}^{+\infty} f(\nu) f^*(\nu) d\nu = \frac{2}{T} \int_0^\infty |f(\nu)|^2 d\nu,\end{aligned}$$

as

$$|f(\nu)|^2 = f(\nu) f(-\nu) = |f(-\nu)|^2.$$

Hence for the total density of radiation we obtain the expression

$$u = \int_0^\infty u_\nu d\nu = \frac{3}{2\pi T} \int_0^\infty |f(\nu)|^2 d\nu,$$

and for the distribution over the spectrum we therefore have

$$u_\nu = \frac{3}{2\pi T} |f(\nu)|^2.$$

After these preliminary remarks on the radiation field we now pass on to the equation giving the vibrations of the linear harmonic oscillator. If the oscillator is capable of vibrating only in the  $x$ -direction, this equation is

$$m\ddot{x} + ax = eE_x(t),$$

and the proper frequency is accordingly given by

$$m(2\pi\nu_0)^2 = a; \quad \nu_0 = \frac{1}{2\pi} \sqrt{\frac{a}{m}}.$$

Now, as we know, the most general solution of a non-homogeneous differential equation is obtained by adding one solution of the non-homogeneous equation to the general solution of the homogeneous equation. This last may be written in the form

$$x(t) = x_0 \sin(2\pi\nu_0 t + \phi),$$

where  $x_0$  and  $\phi$  are two arbitrary constants. The expression

$$x(t) = \frac{e}{2\pi\nu_0 m} \int_0^t E_x(t') \sin 2\pi\nu_0(t - t') dt'$$

is a solution of the non-homogeneous equation and satisfies the initial



conditions  $x(0) = 0$  and  $\dot{x}(0) = 0$ .  $x(0) = 0$  is evident; to prove the other statements we carry out the differentiation:

$$\begin{aligned}\dot{x}(t) &= \frac{e}{2\pi\nu_0 m} [E_x(t') \sin 2\pi\nu_0(t - t')]_{t'=0} \\ &\quad + \frac{e}{m} \int_0^t E_x(t') \cos 2\pi\nu_0(t - t') dt';\end{aligned}$$

here the first term vanishes, and we see that  $\dot{x}(0) = 0$ . Then

$$\begin{aligned}\ddot{x}(t) &= \frac{e}{m} [E_x(t') \cos 2\pi\nu_0(t - t')]_{t'=0} \\ &\quad - \frac{2\pi\nu_0 e}{m} \int_0^t E_x(t') \sin 2\pi\nu_0(t - t') dt' = \frac{e}{m} E_x(t) - \frac{a}{m} x(t);\end{aligned}$$

that is, the expression given above is actually a solution of the non-homogeneous differential equation.

We now proceed to investigate the work done by the field on the oscillator. From the differential equation of the vibrations we readily see (multiplication by  $\dot{x}$  and integration with respect to the time leads to the energy theorem) that the work done per unit time is given by

$$\delta W = \frac{e}{\tau} \int_0^\tau \dot{x}(t) E_x(t) dt.$$

Now it is obvious that the part of the work done which arises from the free vibration (the solution of the homogeneous equation) vanishes. Hence the work done per second is obtained by evaluating the integral of the remaining part only:

$$\delta W = \frac{e}{\tau} \frac{e}{m} \int_0^\tau E_x(t) dt \int_0^t E_x(t') \cos 2\pi\nu_0(t - t') dt'.$$

The integrand is obviously symmetrical in  $t$  and  $t'$ ; hence the expression for  $\delta W$  may be transformed in the following way. We see immediately that

$$\delta W = \frac{e^2}{m\tau} \int_0^\tau E_x(t') dt' \int_{t'}^\tau E_x(t) \cos 2\pi\nu_0(t - t') dt.$$

For in the first instance we have to integrate with respect to  $t'$  between the limits 0 and  $t$  and subsequently with respect to  $t$  between the limits 0 and  $\tau$ ; but of course we obtain the same result by integrating first with respect to  $t$  from  $t'$  to  $\tau$  and then with respect to

$t'$  from 0 to  $\tau$ . If we now merely interchange the letters  $t$  and  $t'$ ,  $\delta W$  may also be written in the form

$$\delta W = \frac{1}{2} \frac{e^2}{m\tau} \int_0^\tau E_x(t) dt \left\{ \int_0^t + \int_t^\tau \right\} E_x(t') \cos 2\pi\nu_0(t - t') dt',$$

or, if we replace  $\cos 2\pi\nu_0(t - t')$  by

$$\frac{1}{2} \{ e^{2\pi i\nu_0(t-t')} + e^{-2\pi i\nu_0(t-t')} \},$$

$$\begin{aligned} \delta W &= \frac{1}{4} \frac{e^2}{m\tau} \left\{ \int_0^\tau E_x(t) e^{2\pi i\nu_0 t} dt \int_0^\tau E_x(t') e^{-2\pi i\nu_0 t'} dt' \right. \\ &\quad \left. + \int_0^\tau E_x(t) e^{-2\pi i\nu_0 t} dt \int_0^\tau E_x(t') e^{2\pi i\nu_0 t'} dt' \right\} \\ &= \frac{e^2}{2m\tau} |f(\nu_0)|^2. \end{aligned}$$

Thus the work done by the field per second on the linear oscillator is

$$\delta W = \frac{e^2}{2m\tau} \frac{2\pi\tau}{3} u_\nu = \frac{\pi e^2}{3m} u_\nu,$$

if we use the formula deduced above for the density of radiation; we thus obtain the expression given in the text.

## XXIX. Temperature and Entropy in Quantum Statistics (p. 211).

The proof of the fact that the quantity  $\beta$  occurring in statistics is inversely proportional to the absolute temperature can be presented in the same form for all three types of statistics, the Boltzmann (B.), the Bose-Einstein (B.E.), and the Fermi-Dirac (F.D.). The fundamental formulæ for the probability of a state for all three statistics can be unified by introducing the symbol

$$\gamma = \begin{cases} 0 & \text{for B.} \\ -1 & \text{for B.E.} \\ +1 & \text{for F.D.} \end{cases}$$

We then have

$$\log W = \sum_s \{ (n_s - \gamma g_s) \log (g_s - \gamma n_s) + \gamma g_s \log g_s - n_s \log n_s \}.$$

The maximum of this, subject to the subsidiary conditions

$$\sum_s n_s = N, \quad \sum_s n_s \epsilon_s = E,$$

is attained for

$$\begin{aligned}\alpha + \beta\epsilon_s &= \frac{\partial \log W_s}{\partial n_s} \\ &= \log(g_s - \gamma n_s) + \frac{\gamma^2 g_s - \gamma n_s}{g_s - \gamma n_s} - \log n_s - 1.\end{aligned}$$

The second term of the last expression is equal to zero for  $\gamma = 0$ , and equal to 1 for  $\gamma = \pm 1$ ; we can therefore replace it simply by  $\gamma^2$ :

$$\alpha + \beta\epsilon_s = \log(g_s - \gamma n_s) - \log n_s + \gamma^2 - 1.$$

Hence

$$n_s = \frac{g_s}{e^{\alpha + \beta\epsilon_s + 1 - \gamma^2} + \gamma}.$$

Here  $\alpha$  and  $\beta$  are to be determined by the two subsidiary condition, (involving  $N$ ,  $E$ ).

Using this value of  $n_s$  in  $\log W$ , we find

$$\log W = \sum_s \{n_s(1 - \gamma^2 + \alpha + \beta\epsilon_s) + \gamma g_s \log(1 + \gamma e^{-\alpha - \beta\epsilon_s})\}.$$

We now assume that the system (a gas) is enclosed in a vessel of variable volume; this volume may be defined, say, by the position co-ordinate  $a$  of a piston. Thus the energy values  $\epsilon_s$  are functions of  $a$ . If changes in  $a$  are made extremely slowly, no quantum jumps are excited by these changes; the numbers  $n_s$  for the quantum states are therefore not changed. Such processes are called "adiabatic" (a better word would be "quasistatic"). The work done in a small change is

$$\delta A = \sum_s n_s \delta \epsilon_s = \sum_s n_s \frac{\partial \epsilon_s}{\partial a} \delta a = -K \delta a,$$

where

$$K = -\sum_s n_s \frac{\partial \epsilon_s}{\partial a}$$

denotes the force (pressure) which opposes the alteration of  $a$ .

If, however, the change of  $a$  takes place rapidly, uncontrollable quantum jumps occur, and therefore alterations in the values  $n_s$ . The corresponding change in the energy is called "heat supplied":

$$\delta Q = \sum_s \epsilon_s \delta n_s.$$

The total change of energy is the sum

$$\delta E = \delta A + \delta Q = \sum_s n_s \delta \epsilon_s + \sum_s \epsilon_s \delta n_s.$$

The change in  $\log W$  is

$$\begin{aligned} \delta \log W = \sum_s n_s \delta(1 - \gamma^2 + \alpha + \beta \epsilon_s) + (1 - \gamma^2 + \alpha) \sum_s \delta n_s \\ + \beta \sum_s \epsilon_s \delta n_s + \gamma \sum_s g_s \delta \log(1 + \gamma e^{-\alpha - \beta \epsilon_s}). \end{aligned}$$

Here it is to be noted that not only  $\epsilon_s$  but also  $\alpha$  and  $\beta$  are to be regarded as functions of  $a$ . Now we have

$$\sum_s \delta n_s = 0,$$

also

$$\gamma g_s \delta \log(1 + \gamma e^{-\alpha - \beta \epsilon_s}) = \frac{\gamma^2 g_s e^{-\alpha - \beta \epsilon_s}}{1 + \gamma e^{-\alpha - \beta \epsilon_s}} \delta(-\alpha - \beta \epsilon_s).$$

For B.E. and F.D.,  $\gamma^2 = 1$ ; hence this last expression is equal to

$$-n_s \delta(\alpha + \beta \epsilon_s).$$

For B.,  $\gamma = 0$ , and, on account of  $n_s = g_s e^{-1 - \alpha - \beta \epsilon_s}$ ,

$$\begin{aligned} \sum_s n_s \delta(1 + \alpha + \beta \epsilon_s) &= \sum_s g_s e^{-1 - \alpha - \beta \epsilon_s} \delta(1 + \alpha + \beta \epsilon_s) \\ &= -\sum_s \delta n_s = 0. \end{aligned}$$

In both cases we are left with

$$\delta \log W = \beta \sum_s \epsilon_s \delta n_s = \beta \delta Q.$$

It follows that  $\beta$  is an integrating factor of the differential of the heat;  $\beta$  must therefore be proportional to  $1/T$ . Putting

$$\beta = \frac{1}{kT},$$

we have

$$\delta Q = kT \delta \log W;$$

and since this, according to the second law of thermodynamics, has to be equal to  $T dS$ , we obtain for the entropy

$$S = k \log W,$$

which is Boltzmann's celebrated formula.

### XXX. Thermionic Emission of Electrons (p. 221).

Here we shall prove the two formulæ which we gave in the text for the thermionic emission of electrons (the Richardson effect) firstly on the basis of classical statistics and secondly on the basis of the

Fermi-Dirac statistics. For this we require to calculate the number of electrons striking one square centimetre of the boundary in the metal per second, such that the kinetic energy of their motion normal to the boundary is sufficient to carry the electron over the energy barrier of height  $\epsilon_a$  which represents the boundary. We have therefore to determine, on the basis of the distribution law, the number of electrons for which, e.g.,

$$\frac{1}{2}mv_x^2 \geq \epsilon_a.$$

We begin with classical statistics. On this basis the number of electrons whose velocity lies between  $v$  and  $v + dv$  is given by

$$dN = 4\pi nV \left( \frac{m}{2\pi kT} \right)^{\frac{3}{2}} e^{-\frac{1}{2}mv^2/kT} v^2 dv$$

(p. 15); similarly (replacing  $4\pi v^2 dv$  by  $dv_x dv_y dv_z$  and integrating) the number of electrons with a velocity component between  $v_x$  and  $v_x + dv_x$  is

$$\begin{aligned} dN_x &= nV \left( \frac{m}{2\pi kT} \right)^{\frac{3}{2}} dv_x \int_{-\infty}^{+\infty} \int_{-\infty}^{+\infty} e^{-\frac{1}{2}m(v_x^2 + v_y^2 + v_z^2)/kT} dv_y dv_z \\ &= nV \sqrt{\frac{m}{2\pi kT}} e^{-\frac{1}{2}mv_x^2/kT} dv_x. \end{aligned}$$

To find the number of electrons falling on unit area of the boundary per second we have first to divide the above number by  $V$ , to get the density of the electrons, and then multiply by  $v_x$ , as in unit time there impinge on the boundary all the molecules with the component  $v_x$  which were contained in the layer of breadth  $v_x$  in front of the boundary (§ 3, p. 5). We thus obtain the emission current by evaluating the integral

$$i = en \sqrt{\frac{m}{2\pi kT}} \int_{\sqrt{(2\epsilon_a/m)}}^{\infty} v_x e^{-\frac{1}{2}mv_x^2/kT} dv_x.$$

This integral may be evaluated exactly, giving

$$i = en \sqrt{\frac{kT}{2\pi m}} e^{-\epsilon_a/kT},$$

which is the expression stated in the text.

The calculation takes a similar course in the case of the Fermi distribution. Here we start from the distribution function (p. 216)

$$dN = \frac{8\pi V}{h^3} \frac{\sqrt{2m^3} \sqrt{\epsilon} d\epsilon}{e^{\alpha + \epsilon/kT} + 1} \quad (\epsilon = \frac{1}{2}mv^2),$$

where  $\alpha$ , the degeneracy parameter, is determined from the subsidiary condition  $\int dN = N$ . If, however, we confine our attention to relatively low temperatures (room temperatures), we can use the approximate formula given on p. 219,

$$dN = \frac{8\pi V}{h^3} \frac{\sqrt{2m^3} \sqrt{\epsilon} d\epsilon}{e^{(\epsilon - \epsilon_0)/kT} + 1} = \frac{8\pi V m^3}{h^3} \frac{v^2 dv}{e^{(\epsilon - \epsilon_0)/kT} + 1},$$

where

$$\epsilon_0 = \frac{h^2}{2m} \left( \frac{3n}{8\pi} \right)^{\frac{2}{3}}$$

represents the zero-point energy. Here, just as before, we obtain the Richardson current by evaluating the integral

$$i = \frac{2m^3 e}{h^3} \int_{-\infty}^{+\infty} \int_{-\infty}^{+\infty} dv_y dv_z \int_{\sqrt{(2\epsilon_a/m)}}^{\infty} \frac{v_x}{e^{(\epsilon - \epsilon_0)/kT} + 1} dv_x.$$

Now in a metal at room temperature ( $\epsilon_a - \epsilon_0$ ) is always very large compared with  $kT$ , ( $\epsilon_a - \epsilon_0$ ) amounting to several electron-volts (p. 222), while  $kT$  at 300° K. corresponds to an energy of about 0.03 electron-volt. Hence in the integrand we always have  $e^{(\epsilon - \epsilon_0)/kT} \gg 1$ , so that we may neglect the 1 in the denominator, thus obtaining the integral

$$i = \frac{2m^3 e}{h^3} e^{\epsilon_a/kT} \int_{-\infty}^{+\infty} \int_{-\infty}^{+\infty} dv_y dv_z \int_{\sqrt{(2\epsilon_a/m)}}^{\infty} v_x e^{-\frac{1}{2}m(v_x^2 + v_y^2 + v_z^2)/kT} dv_x.$$

The integrations with respect to  $v_y$  and  $v_z$  are equivalent to evaluations of Gaussian error integrals (see Appendix I, p. 259); the integration with respect to  $v_x$  can, as before, be carried out by elementary means, giving

$$i = \frac{4\pi em}{h^3} (kT)^2 e^{-(\epsilon_a - \epsilon_0)/kT},$$

the law stated in the text.

### XXXI. Theory of Valency Binding (p. 253).

The fundamental idea of the Heitler-London theory of valency binding is as follows. As a model of the hydrogen molecule we imagine two nuclei  $a$  and  $b$  on the  $x$ -axis at a distance  $R$  apart, and two electrons 1 and 2 revolving about the nuclei. To the state of two widely-separated neutral atoms there corresponds a large value of  $R$  and a motion of the electrons such that each one revolves round one of the two nuclei. Let the two atoms be in the ground state and have the

same energy  $E_0^1 = E_0^2 = E_0$ . The motion of the electrons is described by proper functions  $u$ , which, relative to the corresponding nuclei, are identical; that is, one is obtained from the other by substituting  $x + R$  for  $x$ ; we shall briefly write them in the form

$$\begin{aligned}\psi_a^{(1)} &= u(x_1, y_1, z_1), \\ \psi_b^{(2)} &= u(x_2 + R, y_2, z_2).\end{aligned}$$

The functions  $u$  are the same as the proper functions of atomic hydrogen (Appendix XVIII (p. 298)). Hence the two Schrödinger equations

$$\begin{aligned}H_a^0 \psi_a^{(1)} &= E_0 \psi_a^{(1)}, \\ H_b^0 \psi_b^{(2)} &= E_0 \psi_b^{(2)},\end{aligned}$$

where  $H^0$  denotes the energy operator of the hydrogen atom, are satisfied identically and the suffixes  $a$  and  $b$  indicate that in the one case the co-ordinates of the electron are relative to the nucleus  $a$ , in the other to the nucleus  $b$  (see above).

The energy operator of the molecule which arises when the atoms approach one another (when  $R$  diminishes) differs from the sum  $H_a^0 + H_b^0$  by the interaction energy of the two atoms,

$$V = e^2 \left( \frac{1}{r_{ab}} + \frac{1}{r_{12}} - \frac{1}{r_{a2}} - \frac{1}{r_{b1}} \right),$$

where  $r_{ab}$  denotes the distance between the two nuclei ( $R$ ),  $r_{12}$  the distance between the two electrons, and  $r_{a2}$ ,  $r_{b1}$  the distances between either electron and the nucleus of the other atom. For the molecule we accordingly have the Schrödinger equation

$$(H_a^0 + H_b^0 + V)\psi^{(1, 2)} = E\psi^{(1, 2)}.$$

We now seek to find an approximate solution for this equation by assuming that (to a first approximation) the function  $\psi^{(1, 2)}$  of the co-ordinates of the two electrons is a product of some proper function  $\psi_a^{(1)}$  of one electron and some proper function  $\psi_b^{(2)}$  of the other electron. Here, however, we have to bear in mind that the state of the system is a degenerate one. The total energy of the two separated atoms,

$$E = E_0^1 + E_0^2 = 2E_0,$$

corresponds not only to the product  $\psi_a^{(1)}\psi_b^{(2)}$  but also to  $\psi_a^{(2)}\psi_b^{(1)}$  and every possible combination of the two expressions. These two vibrational forms will interact much more strongly with each other than with the vibrations corresponding to any other energy levels, owing

to the coupling as the two atoms approach one another. They accordingly suffice as a basis for a rough approximation, i.e. we attempt to represent the function  $\psi^{(1, 2)}$  approximately by a linear combination of the two functions  $\psi_a^{(1)}\psi_b^{(2)}$  and  $\psi_a^{(2)}\psi_b^{(1)}$ . Instead of these we may also start from the symmetric and antisymmetric combinations

$$\begin{aligned}\psi_{\text{sym.}} &= \psi_a^{(1)}\psi_b^{(2)} + \psi_a^{(2)}\psi_b^{(1)}, \\ \psi_{\text{antis.}} &= \psi_a^{(1)}\psi_b^{(2)} - \psi_a^{(2)}\psi_b^{(1)},\end{aligned}$$

a course which has two advantages: (1) closer investigation shows that to a first approximation the symmetric function and the antisymmetric function are not coupled with one another by the Schrödinger equation, i.e. that each function by itself represents a separate state of the molecule; (2) they may readily be distinguished by means of the spin. For according to the exclusion principle the proper functions of a system must be antisymmetric in all the co-ordinates of the two electrons (of course taking the spin into account; cf. p. 209). If we were to give the electrons spin variables, as we did in the case of atomic spectra (§ 8, p. 169), the corresponding spin function would have to be antisymmetric for  $\psi_{\text{sym.}}$  and symmetric for  $\psi_{\text{antis.}}$  in order to satisfy the exclusion principle. This means that in the case of  $\psi_{\text{sym.}}$  the spins are antiparallel and balance one another, whereas in the case of  $\psi_{\text{antis.}}$  they are parallel and additive.

Now the perturbation calculation shows that when the two atoms are brought closer together (when the coupling is increased) the proper value of the uncoupled system ( $2E_0$ ) is split up into two values

$$E_1 = 2E_0 - W_1, \quad E_2 = 2E_0 + W_2,$$

where the functions  $W(R)$  have the following meanings:

$$W_1 = \frac{H_1 S + H_2}{1 + S}, \quad W_2 = \frac{H_1 S - H_2}{1 - S},$$

$$H_1 = \iint (\psi_a^{(1)})^2 (\psi_b^{(2)})^2 V d\tau_1 d\tau_2,$$

$$H_2 = \iint \psi_a^{(1)} \psi_a^{(2)} \psi_b^{(1)} \psi_b^{(2)} V d\tau_1 d\tau_2,$$

$$S = \iint \psi_a^{(1)} \psi_a^{(2)} \psi_b^{(1)} \psi_b^{(2)} d\tau_1 d\tau_2,$$

the integrations being taken over the co-ordinates of the two electrons. As  $\psi^2$ , apart from a factor, represents the density of the charge-cloud of the electron, the first integral represents the Coulomb force due to the mutual actions of the charges distributed over each



atom. The second integral is characteristic of the quantum theory, as here the proper functions are not related quadratically, so that an interpretation based on charge-densities is not possible. This integral is known as the "exchange integral".

The evaluation of the integral as a function of  $R$ , which is very troublesome, gives the curves shown in fig. 7, p. 243 for the energy as a function of the distance between the two atoms. The function  $\psi_{\text{sym.}}$ , which is symmetrical in the co-ordinates of the electrons, gives the lower curve, according to which the energy for a given intramolecular distance has a minimum, while the antisymmetric function  $\psi_{\text{antis.}}$  gives the monotonically ascending branch of the curve, which corresponds to a permanent repulsion between the two atoms. It turns out, therefore, that the state that leads to binding is that for which the electronic spins balance one another. Thus it comes about that we may regard the spin as a physical substitute for the chemical valency. But, as was mentioned in the text, in the case of polyatomic molecules this idea leads to difficulties which have only partly been overcome (Heitler, Rumer, Weyl).

### XXXII. Theory of the van der Waals Forces (p. 291).

According to London, the theory of the van der Waals forces rests on a fact which is a distinctive feature of the quantum theory, namely, the existence of a finite zero-point energy (cf. the case of the harmonic oscillator (Appendixes XV, p. 291, XVI, p. 295)). According to the classical theory, the state of least energy of an oscillator is that of zero energy; this is the state of rest in the position of equilibrium. According to wave mechanics, however, the ground state has a finite energy  $E_0 = \frac{1}{2}\hbar\nu_0$ , and the corresponding proper function is the Gaussian error function  $\psi_0 = ae^{-\frac{1}{2}aq^2}$ , where  $a = \frac{4\pi^2 m \nu_0}{h}$  and  $\nu_0$  is the proper frequency of the oscillator. This zero-point energy can be explained by Heisenberg's uncertainty principle, according to which if the energy of a particle is prescribed accurately its position cannot be accurately determined. The proper function, which is a Gaussian error function, expresses the uncertainty directly. For the error curve immediately gives the mean square deviation of the co-ordinate (as  $\bar{q} = 0$ ):

$$\overline{\delta q^2} = \overline{q^2} = \frac{\int_{-\infty}^{+\infty} q^2 \psi_0^2 dq}{\int_{-\infty}^{+\infty} \psi_0^2 dq} = \frac{1}{2a} = \frac{h}{8\pi^2 m \nu_0}.$$

On the other hand, from the energy equation we have

$$E = \frac{1}{2}p^2/m + \frac{1}{2}m(2\pi\nu_0)^2q^2;$$

if the energy is accurately determined, the mean square deviation of the momentum is

$$\overline{\delta p^2} = m^2(2\pi\nu_0)^2\overline{\delta q^2} = \frac{1}{2}h\nu_0m.$$

Hence

$$\overline{\delta q^2}\overline{\delta p^2} = \frac{h^2}{16\pi^2}$$

or, with  $(\Delta q)^2 = \overline{\delta q^2}$ ,  $(\Delta p)^2 = \overline{\delta p^2}$ :

$$\delta p \delta q = \frac{h}{4\pi}.$$

This is the exact form of Heisenberg's uncertainty principle  $\delta p \delta q$  (see Appendix XII, p. 282, and Appendix XXII, p. 312).

After this digression on the zero-point energy and the theoretical uncertainties of the position and the momentum when the energy is determined accurately, we now return to London's explanation of the occurrence of the van der Waals forces. As a simple model we consider two linear oscillators at a distance  $R$  apart, vibrating in the direction of the line joining them (the  $x$ -axis). We think of these oscillators as vibrating electrical dipoles in which the positive charges  $e$  are held fast in the position of equilibrium while the negative charges  $-e$  vibrate about these equilibrium positions, their displacements being  $x_1$  and  $x_2$ . We express the restoring forces on the oscillators in the form  $-\frac{e^2}{a}x_1$  and  $-\frac{e^2}{a}x_2$ ; their potential energies are then  $\frac{1}{2}\frac{e^2}{a}x_1^2$  and  $\frac{1}{2}\frac{e^2}{a}x_2^2$ . In addition there is the coupling force acting between the two oscillators, for which we assume Coulomb's law of force. The potential energy of this interaction is

$$\frac{e^2}{R} + \frac{e^2}{R+x_1+x_2} - \frac{e^2}{R+x_1} - \frac{e^2}{R+x_2},$$

or, if we assume that  $R$  is very large compared with  $x_1$  and  $x_2$ , and expand, we have

$$\frac{2e^2x_1x_2}{R^3}.$$

Hence the energy equation for the two oscillators is of the form

$$W = \frac{1}{2m} (p_1^2 + p_2^2) + \frac{e^2}{2a} (x_1^2 + x_2^2) + \frac{2e^2 x_1 x_2}{R^3}.$$

In the absence of the coupling term the resonators would themselves vibrate with the frequency

$$\nu_0 = \frac{1}{2\pi} \sqrt{\frac{e^2}{am}}.$$

If the coupling is taken into account the frequency, as we have repeatedly stated before, is split up. Here we shall carry out the actual calculation of the splitting. For this purpose it is most convenient to refer the quadratic expression for the potential energy to "principal axes"; this does not affect the form of the kinetic energy. This is done by means of the transformation

$$x_s = \frac{1}{\sqrt{2}} (x_1 + x_2), \quad x_a = \frac{1}{\sqrt{2}} (x_1 - x_2);$$

as  $p = m\dot{x}$ , we also have

$$p_s = \frac{1}{\sqrt{2}} (p_1 + p_2), \quad p_a = \frac{1}{\sqrt{2}} (p_1 - p_2).$$

If we substitute these new co-ordinates and momenta in the expression for the energy, we have

$$W = \frac{1}{2m} (p_s^2 + p_a^2) + \frac{e^2}{2a} (x_s^2 + x_a^2) + \frac{e^2}{R^3} (x_s^2 - x_a^2),$$

or, in another form,

$$W = \left\{ \frac{1}{2m} p_s^2 + \left( \frac{e^2}{2a} + \frac{e^2}{R^3} \right) x_s^2 \right\} + \left\{ \frac{1}{2m} p_a^2 + \left( \frac{e^2}{2a} - \frac{e^2}{R^3} \right) x_a^2 \right\}.$$

This, however, is the energy equation for two non-coupled oscillators vibrating with the two frequencies

$$\nu_s = \frac{1}{2\pi} \sqrt{\frac{e^2}{m} \left( \frac{1}{a} + \frac{2}{R^3} \right)}, \quad \nu_a = \frac{1}{2\pi} \sqrt{\frac{e^2}{m} \left( \frac{1}{a} - \frac{2}{R^3} \right)}.$$

Hence the quantised energy of the system is

$$E_{n_s n_a} = h\nu_s(n_s + \frac{1}{2}) + h\nu_a(n_a + \frac{1}{2}),$$

which depends on  $R$ , as the new frequencies are functions of the

distance between the oscillators. For the ground state we obtain the zero-point energy of the two vibrations of the oscillator,

$$E_{00} = \frac{1}{2}h(\nu_s + \nu_a),$$

that is, if we again expand,

$$\begin{aligned} E_{00} &= \frac{h}{4\pi} \left\{ \sqrt{\frac{e^2}{m} \left( \frac{1}{a} + \frac{2}{R^3} \right)} + \sqrt{\frac{e^2}{m} \left( \frac{1}{a} - \frac{2}{R^3} \right)} \right\} \\ &= \frac{h}{2\pi} \sqrt{\frac{e^2}{ma}} \left( 1 - \frac{a^2}{2R^6} + \dots \right) = h\nu_0 \left( 1 - \frac{a^2}{2R^6} + \dots \right). \end{aligned}$$

The additional energy is therefore negative and inversely proportional to the sixth power of the distance between the oscillators; the oscillators will accordingly attract one another with a force varying as the inverse seventh power of the distance between them. The magnitude of the attraction also depends on  $\nu_0$  and on the square of the constant  $a$ , which is obviously a measure of the deformability of the oscillators.

Exactly similar considerations apply to any atomic system in which there is interaction between the constituents, and invariably lead to the result that between two systems in the ground state there will be a force of attraction whose potential energy is inversely proportional to the sixth power of the distance and whose magnitude is proportional to the product of the deformabilities of the two atoms.

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tion thus occurs before the convergence limit of the bands, i.e., by longer wave light than corresponds to the dissociation continuum.

As in the case of the phenomenon of preionization discussed on page 186, predissociation is caused by nonradiative transitions from quantized energy states into a region of continuous energy, which in the case of predissociation corresponds to that of the dissociated molecule. To explain predissociation, Fig. 182 shows the potential curves and vibration level diagrams of two excited molecular states. We assume that a transition from a lower state, not shown, to the state  $a'$  is possible.

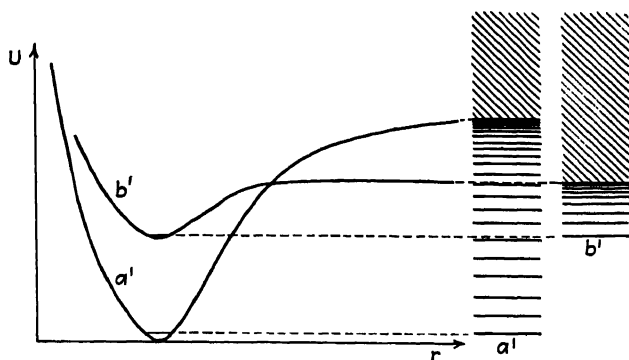


FIG. 182. Potential curves and corresponding vibration term sequences leading to predissociation of a diatomic molecule.

Then in the absence of state  $b'$  we would observe a sequence of absorption bands which corresponds to transitions from a lower state to all  $a'$  states beyond which, after the dissociation limit of  $a'$  is reached, the dissociation continuum of  $a'$  follows. Now we consider the influence of the disturbing molecular state  $b'$ , whose dissociation energy lies below that of  $a'$ . If certain conditions (selection rules) between the electron quantum numbers of  $a'$  and  $b'$  are fulfilled, there exists for the discrete  $a'$  levels lying above the dissociation limit of  $b'$  the possibility of non-radiative transitions (without gain or loss of energy) into the continuous energy region of  $b'$ , and thus of dissociation before the dissociation limit of  $a'$  itself is reached. Since the Franck-Condon principle is valid also for transition without emission or absorption of radiation, these transitions from  $a'$  into the continuum of  $b'$  occur with high probability only in the region of the intersection of  $a'$  and  $b'$ , while for higher states the probability of predissociation again decreases.

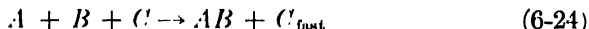
As in the case of preionization (page 186), the possibility of transitions into the continuum implies a considerable reduction of the life-time of the discrete states of  $a'$  and thus, according to the uncertainty rela-

tion, an increase in the width of these energy states and the lines resulting from their combination. In the region of predissociation this line width usually becomes so large that the different rotation lines of a band merge into one another, thus accounting for the diffuse appearance of the dissociation bands. Predissociation is observed only in absorption, because, as a result of the nonradiative transitions into the continuous energy region, the radiation probability of predissociating states is extremely small.

For the photochemist this section contains some very important results: *Irradiation by those wavelengths which correspond to sharp absorption bands of the irradiated molecule produce excited, and thus reactive, molecules. Irradiation by those wavelengths, on the other hand, which correspond to continuous absorption spectra or diffuse absorption bands (which appear diffuse only because of the small dispersion of the spectroscopist), produces dissociation of the irradiated molecules.* Thus the relation between the continuous absorption spectra of molecules and their decomposition into atoms is explained.

### c. The Processes of Formation of Molecules from Atoms

What do we know about the reverse process, the formation of (diatomic) molecules from their atoms? In this case the situation is similar to the combination of an ion and an electron to form a neutral atom as discussed on page 105. If two atoms which can form a molecule do collide, they form a molecule only if the binding energy released in the collision (which of course is numerically equal to the dissociation energy) is somehow carried away. This can be done by a third atom. In that case we have the formation of a molecule in a triple collision of atoms and can write the process as an equation,



We now consider whether this molecular recombination can also occur if *two* atoms collide and the binding energy is emitted as a photon. In order to understand this problem correctly, we must first know the difference between atomic (homopolar) and ionic (heteropolar) molecules. By *atomic molecules* we mean those which decompose into normal atoms when dissociated in the ground state (e.g.,  $\text{H}_2$ ,  $\text{O}_2$ ,  $\text{Cl}_2$ , etc.). *Ionic* or *heteropolar molecules*, on the other hand, are those molecules which (like  $\text{NaCl}$ ) dissociate from the ground state into ions ( $\text{Na}^+$  and  $\text{Cl}^-$ ).

In the usual case of atomic molecules the colliding atoms  $A$  and  $B$  "run along" the potential curve of the normal molecule (see Fig. 180). *The formation of a molecule, under emission of the binding energy and the*

relative kinetic energies of  $A$  and  $B$ , is in this case just as impossible as the reverse process, the optical dissociation without electron excitation. The reason is that transitions from the continuous energy region to the vibrational ground state of the same electron state do not occur, because their transition probability, according to the Franck-Condon principle, is practically zero (Fig. 180). Two normal atoms which can form a molecule thus cannot recombine in a two-particle collision, but only in a triple collision in which the third partner, which can be the wall of a glass tube, for example, takes over the excess energy.

The situation is different in the case of an ionic molecule (e.g.,  $\text{NaCl}$ ).

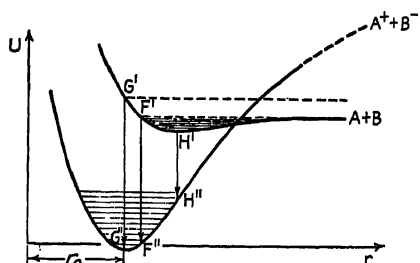
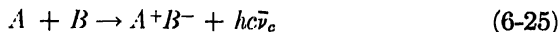


FIG. 183. Set of potential curves for explaining the formation of an ionic molecule in a collision of two atoms upon emission of the recombination continuum.

Here the colliding atoms are in a different electron state (curve  $V'$  of Fig. 183) than the ground state  $V''$  which is formed by two ions. Consequently, a recombination in a two-particle collision is possible. In this case, the radiation probability (and with it the recombination probability) is determined by the probability of the electron transition  $AB \rightarrow A^+B^-$  and by the Franck-Condon principle. In the case of the alkali halides, for ex-

ample, transitions to the normal ground state can occur, according to the Franck-Condon principle, from the reversal points between  $F'$  and  $G'$  of the colliding normal atoms  $A$  and  $B$ . In this case the energy is radiated away as an emission continuum whose wavelength limits are given by the transition arrows  $F'F''$  and  $G'G''$ , so that we have the two-particle collision recombination process



Consequently, just as photodissociation under continuous absorption is possible only if there is a simultaneous change of the electron configuration, the formation of molecules in a two-particle collision under continuous emission of the binding and kinetic energies occurs only if there is a simultaneous change in the electron configuration. Such an electron transition is also connected with the radiative recombination of normal and excited atoms in two-particle collisions, according to the equation



In a highly dissociated and excited gas, for example in a discharge plasma, collisions between normal and excited atoms are sufficiently frequent. These take place "along" the curve of the excited molecular

state, Fig. 181. As in the case mentioned above, in collisions the kinetic energy of the colliding atoms is converted into potential energy so that a point above the dissociation limit on the left side of the potential curve is reached. By transitions from this branch of the potential curve of the excited molecule to the ground state, the colliding atoms  $A$  and  $B^*$  recombine under emission of continuous radiation.

Thus, the observation of such emission continua of molecules furnishes evidence for the process of radiative recombination, and from it the probability can be determined as a function of the velocity of the colliding atoms. Emission continua, which correspond to the recombination process (6-25), are found especially in the case of alkali-halide vapors, whereas continua corresponding to radiative recombination (6-26) have been observed for the halogens and tellurium.

### 6-7. Limitations of the Molecule Concept. Van der Waals Molecules and Collision Pairs

We have so far used the term "molecule" as if it were a self-evident concept, but now we must refine our knowledge and proceed to a higher point of view. By a molecule we usually mean, without giving it too much consideration, a system of two or more atoms or groups of atoms whose potential energy for a definite nuclear configuration is at a minimum, and which is, to a certain extent, independent of its environment. The latter means that (at least at the instant of observation) the interaction within the system is great compared to that between the system and its environment. Physically such a molecule is, as we know, characterized by the forces acting between the atoms, i.e., by the change of potential with the internuclear distance (potential curve) and by the value of its total energy. The energy is measured upward (positive) from the ground state of the normal molecule or (as we do here) from the zero point of the free atoms. In this case the binding energy of the stable molecular states is negative, whereas the kinetic energy of the free, dissociated states of the molecule is positive.

We now can distinguish between three types of molecules, depending on the shape of their potential curves, i.e., their binding forces (lower part of Fig. 184). We consider first the normal, strongly bound molecules which were almost exclusively treated in the last sections. Their characteristic features are a pronounced potential minimum at a relatively small internuclear distance ( $r_0 \sim 1\text{\AA}$ ), and a relatively large dissociation energy of 1 to 5 eV (Fig. 184, curve *a*). The second type of molecule is called a "van der Waals molecule" for reasons which will become clear below. It is characterized by the potential curve *b*, with a shallow minimum at large internuclear distance ( $r_0 \sim 3$  to  $5\text{\AA}$ ), and by a usually

very small dissociation energy (order of magnitude of 0.01 to 0.1 ev in the ground state). These molecules consist of atoms which are not able to form a normal, strongly bound molecule by rearranging their electron shells. However, there exists between them some attraction due to mutual polarization and related interaction forces of the second order. These second-order forces imply a possibility of formation of loosely bound molecules. The same types of interatomic and intermolecular attractive forces which produce these weak bonds are also responsible

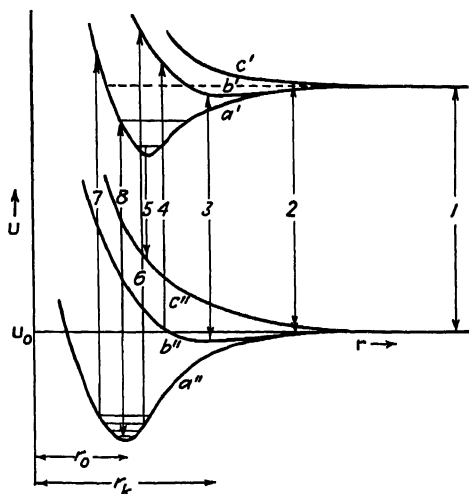


Fig. 184. Potential-curve diagram for the electron ground state and an excited electron state of a diatomic molecule with transitions indicating all possible discrete and continuous molecular spectra.

for the deviations in the behavior of a real from an ideal gas. Since these deviations are described by the van der Waals equation of state, we call those molecules which are bound by these forces van der Waals molecules. Now if we consider as the third type of molecule an atomic system which is characterized by the potential curve  $c$ , we see that this differs from a van der Waals molecule only in that it cannot have negative energy values at all. The repulsive forces always exceed the attractive forces; in every other respect such a system of atoms behaves just as an arbitrary real molecule in a free state of positive energy. It is represented by two atoms in the instant of the collision, if only repulsive forces act between them. Such a system is called a *collision pair*. We introduce it here as the limiting case of the molecule, although it lacks a potential minimum and thus does not fit the customary definition of a molecule given above. Evidently a sharp line cannot be drawn between the van der Waals molecule and the collision pair. Furthermore,



a collision pair can emit and absorb spectra just as other molecules. Finally, a pure repulsive potential curve characterizing a collision pair occasionally also occurs in the case of excited states of regular molecules, as we know, for instance, from Fig. 119.

The best known example of a van der Waals type molecule is the

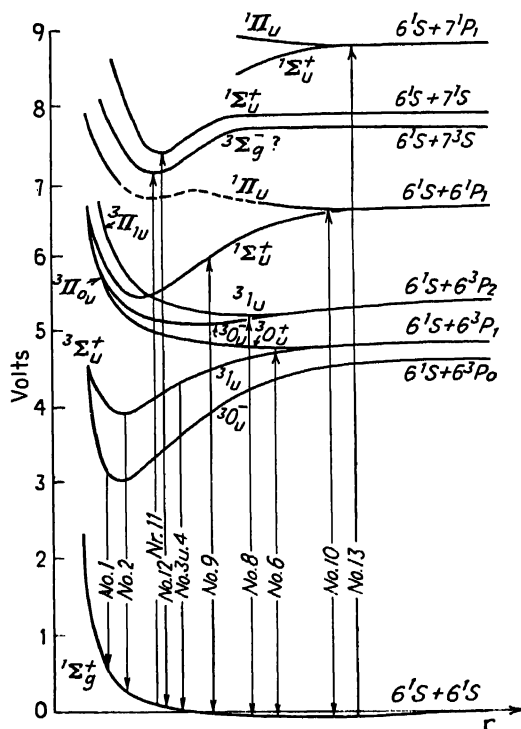


FIG. 185. Potential-curve diagram of the  $\text{Hg}_2$  van der Waals molecule, determined from the study of the diffuse bands and continua. (By Mrozowski and the author.)

$\text{Hg}_2$  molecule consisting of two Hg atoms. All types of binding which we have mentioned above occur among its excited energy states. Figure 185 shows its potential curve diagram which is the product of many years of patient research of a large number of investigators. The  $\text{He}_2$  molecule consisting of two helium atoms, whose existence was established spectroscopically, is another example of a van der Waals molecule, but only in the ground state. Pronounced potential minima are characteristic for the excited states of  $\text{He}_2$  and give rise to the emission of numerous  $\text{He}_2$  bands, which were analyzed, especially by Weizel. By transitions from these stable excited states to the unstable ground state, an emission continuum lying in the far ultraviolet is emitted.

A noticeable van der Waals attraction also exists between two  $O_2$  molecules and, according to an investigation of the author, is responsible for the existence of the double molecule  $O_2-O_2 = (O_2)_2$ . It consists of two normal oxygen molecules and, besides other spectra, absorbs the continuous bands which are responsible for the blue color of compressed as well as liquid oxygen.

Van der Waals forces, whose potential corresponds to curve *b* in Fig. 184, play an important role in other phenomena also. We mention here only the large number of cohesion and adhesion phenomena which are produced by them, as well as the secondary forces by which smaller and larger molecules are bound together in certain crystals and colloidal micelles which play an important role in biology. They also are, at least partly, responsible for the process of polymerization of gigantic molecular complexes from single molecules (e.g., the caoutchouc molecule from about 2,000 isoprene molecules). We point out, finally, that the loose binding of two molecules effected by such van der Waals forces seems to be important for the initial stage of many chemical reactions, because it is often a prerequisite for a sufficiently intense interaction of the molecules. The often very casual treatment of the van der Waals forces and the spectra resulting from them (also the variation of an atomic spectrum by the van der Waals binding of the atom to a large complex, such as a crystal or a wall) does not seem to be justified, since a careful investigation of these phenomena may produce many interesting results.

After this extension of our knowledge about the concept of molecules and the potential curves of the different types of molecules, we return to the discrete as well as continuous spectra which result from the shape and mutual positions of the potential curves of the combining molecular states and, conversely, can be used to determine them. For this purpose we consider the schematic diagram Fig. 184 in which are shown as asymptotes to the right the two families of curves of the ground state and an excited state of an atom. The potential curves originate from bringing an additional normal atom into interaction with the first one, this being in its different states. The different possible types of transitions are represented by arrows. However, we have to consider that the transitions naturally occur also from all neighboring points of the potential curve under consideration.

Whereas the transition number 1 between the undisturbed atomic states produces a sharp spectral line, a broadened atomic line results from transition number 2. It is evident that the amount and character of the broadening depends upon the average interatomic distance (function of the density of the gas), on the temperature, and the shape of the two potential curves. With these few remarks, we have consid-

ered the process of collision broadening, already discussed on page 187, from an entirely new point of view. We considered the emitting or absorbing and the colliding atom together as a molecule (or collision pair) and thus conceived of collision broadening as a limiting case of a molecular spectrum. It is now evident that it is only a small step from the simple broadened atomic line to the spectra of the weakly bound van der Waals molecule represented by transitions 3 and 4. They consist of narrow continuous or apparently continuous bands which are closely related to the corresponding atomic line (transition number 1). The emission and absorption of mercury vapor at pressures which are not too low provide numerous examples of this type of band associated with atomic lines. We have already referred to the identification of these bands as spectra of a  $\text{Hg}_2$  van der Waals molecule. The transitions 5 to 8 occur in the region of nuclear separation in which the electron clouds of the two atoms penetrate each other considerably and thus cause major changes of the energy compared to the unperturbed atomic states. They represent, therefore, molecular spectra in the more restricted and customary sense. Whereas the normal electron band spectra correspond to transitions 8 between the discrete states of both potential curves, there exists an extended emission continuum due to transition 5 from the minimum of an excited molecular state to the repulsion curve of a lower state. The best known examples of this are the continuum of the helium molecule in the extreme ultraviolet, which was mentioned on page 77, and the hydrogen molecular continuum extending from the green to the far ultraviolet. The latter is emitted by any glow discharge in dry hydrogen at several millimeters pressure. In conclusion, transition 7 from the discrete states of the lower potential curve to the continuous energy region of an upper curve accounts for an extended absorption continuum. The best known examples for this are the absorption continua of the halogen molecules  $\text{I}_2$ ,  $\text{Br}_2$ ,  $\text{Cl}_2$ ,  $\text{IBr}$ , etc., which lie in the visible or near ultraviolet.

This discussion reveals how many important conclusions with respect to emission and absorption spectra can be drawn from the potential curve diagram. It is worth while, therefore, to become thoroughly familiar with the concept of the potential curve and its consequences.

#### **6-8. Molecular Rotation and the Determination of Moments of Inertia and Internuclear Distances from the Rotation Structure of the Spectra of Diatomic Molecules**

After we have become acquainted with the numerous phenomena which are related to molecular vibration and dissociation, we shall discuss now molecular rotation and the conclusions which may be

drawn from investigations of the rotational phenomena in molecular spectra.

### a. Rotation Term Diagrams and Infrared Rotation Spectra

The simplest model of the rigid molecule rotating about the axis of greatest moment of inertia, the rotator with free axis, was discussed wave-mechanically on page 218. We found there that only the following sequence of discrete energy values is possible for the rotator:

$$E_{\text{rot}} = hcBJ(J + 1) \quad J = 0, 1, 2, \dots \quad (6-27)$$

where

$$B = \frac{h}{8\pi^2 cI} \quad (6-28)$$

is the rotation constant and

$$I = \mu r^2 = \frac{m_1 m_2}{m_1 + m_2} r^2 \quad (6-29)$$

again is the moment of inertia of the molecule about the axis of rotation.  $J$  is the rotation quantum number belonging to the total angular momentum  $\mathbf{J}$  of the rotating molecule.

The rotation term diagram shown in Fig. 186 follows from (6-27). For optical transitions between these rotation states the selection rule

$$\Delta J = \pm 1 \quad (6-30)$$

FIG. 186. Possible transitions in the rotation-level diagram of the electron and vibration ground state of a diatomic molecule (rotation band).

which follows from the correspondence principle (page 164) or from wave mechanics, applies. Thus each rotation state can combine under emission or absorption of radiation only with its two neighboring states. From this selection rule and Fig. 186 it follows that the rotation spectrum of a diatomic molecule lying in the far infrared consists of a sequence of equidistant spectral lines (rotation lines) with the wave numbers  $2B, 4B, 6B, \dots$ . The distance between two consecutive lines thus should be  $2B$ . Pure rotation spectra in good agreement with this theory were first found by Czerny in the case of the HCl molecule. Just as in the case of the vibration spectrum (page 369), the rotation spectrum can be observed only if a change of the electric moment is associated with the rotation (dipole radiation). Thus the centers of the electric charges must not coincide with the center of mass. This condition is fulfilled only for the unsymmetrical molecules such as HI, but not for  $\text{H}_2$ ,  $\text{Cl}_2$ ,  $\text{O}_2$ , etc. Consequently, the latter do not have easily observable rotation or rotation vibration bands.

If the rotation spectrum is observable, then the rotation constant  $B$  can be taken from it immediately. The moment of inertia of the molecule, and thus one of the most important molecular constants, can be computed from  $B$  according to Eq. (6-28). Moreover, if the masses  $m_1$  and  $m_2$  of the constituent atoms of a molecule are known (and this will always be the case if we know the molecule responsible for the spectrum), then the internuclear distance of the molecule follows immediately from the moment of inertia by means of Eq. (6-29). The internuclear distances determined in this way are within the order of magnitude of 1 Å, which is in agreement with our knowledge of the dimensions of atoms and molecules. Table 14 presents the most important data for a number of diatomic molecules, as determined from their spectra. If the rotation and vibration are optically inactive, as in the case of molecules consisting of identical atoms, then  $B$ , and with it  $I$  and  $r_0$ , can be determined, though not quite as easily, from the rotational structure of the electronic bands, or from the rotational Raman effect, page 352.

Table 14. *Internuclear Distances, Moments of Inertia, Fundamental Vibrational Quanta, and Dissociation Energies of Some Important Molecules in the Ground State*

Molecule	Internuclear distance, $r_0$ Å	Moment of inertia $I$ , $10^{-40}$ gm cm <sup>2</sup>	Fundamental vibrational quantum, cm <sup>-1</sup>	Dissociation energy	
				Electron volt	kcal/mole
H <sub>2</sub>	0.77	0.47	4390	4.46	103
O <sub>2</sub>	1.20	19.1	1580	5.09	118
N <sub>2</sub>	1.09	13.8	2360	7.35	170
S <sub>2</sub>	1.60	67.7	727	4.45	103
Cl <sub>2</sub>	1.98	113.5	565	2.47	57
Br <sub>2</sub>	2.28	342	324	1.96	45
I <sub>2</sub>	2.66	741	214	1.53	35
CO	1.13	15.0	2169	9.67	2207
NO	1.15	16.3	1907	5.3	122
HCl	1.27	2.60	2989	4.40	102

### b. The Rotation Vibration Spectrum

From the discussion of the pure rotation spectrum we proceed to the rotation vibration spectrum which is produced by changes in the rotational and vibrational states of the molecule (without a change in the configuration of the electrons). Two sequences of rotational terms, belonging to the vibration levels  $v = 0$  and  $v = 1$ , are shown in Fig. 187. Corresponding to the selection rule (6-30), the resulting rotation vibration band consists of two "branches" which belong to the transitions  $\Delta J = +1$  and  $\Delta J = -1$ . They are shown to the left and right in

Fig. 187. The line sequence belonging to  $\Delta J = +1$  extends from the zero line (not shown) toward shorter wavelengths and is called the positive or *R* branch. The sequence belonging to  $\Delta J = -1$  extends toward the longer wavelengths and is called the negative or *P* branch.

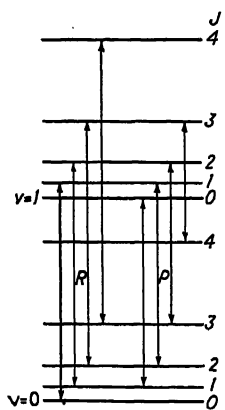


Fig. 187. Energy-level diagram for explaining a rotation vibration band of a diatomic molecule.

If we assume as a rough approximation that the moment of inertia of the molecule is the same for the upper as for the lower vibrational state in spite of the difference in vibration, then the separations of the rotation levels are equal in both vibration states. This case is presented in Fig. 187. Then once more we have equidistant lines separated by the constant distance  $2B$  (Fig. 188). The zero line, which cannot appear because the transition  $\Delta J = 0$  is forbidden, can be easily determined. Its frequency gives the energy of the pure vibrational transition. For every vibrational transition which is possible according to Figs. 177 and 178, we expect a band of the kind described. This corresponds roughly to the spectroscopic results. A small deviation from the assumed constancy of the distances between consecutive lines, however,

is evidence that, first, the moment of inertia is not constant and independent of the vibrational state and that, second, the rotation and vibration of the molecule are not independent. The energy values of the pure rotation and vibration thus cannot simply be added, but the interaction of vibration and rotation must be taken into account by introducing a mixed term containing  $v$  and  $J$  in the

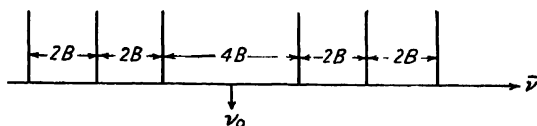


Fig. 188. Schematic representation of a rotation vibration band.

expression for the molecular energy. If this is done, complete agreement between theory and spectroscopic results is achieved.

### c. The Rotation Structure of Normal Electron Bands

We now take up the most complicated case, that of band spectra resulting from simultaneous changes in the electron configuration, vibration, and rotation of a molecule. These bands occur in the visible and ultraviolet spectral regions. As a result of the participation of an

electron transition in the absorption and emission, there is always a change in the electric moment of the molecule. Electron band spectra thus are emitted and absorbed also by symmetrical molecules such as  $\text{H}_2$ ,  $\text{O}_2$ , etc. These electron transitions also make it possible for transitions (except for the  $\Sigma \rightarrow \Sigma$  transitions, see below) to occur without a change in the rotational quantum number. Thus, instead of (6-30), the rotational selection rule for electron band spectra is

$$\Delta J = 0 \text{ or } \pm 1 \quad (6-31)$$

with the limitation that transitions between states with zero rotation ( $0 \leftrightarrow 0$ ) are forbidden since there is no change of the electric moment connected with them. Consequently, a zero or  $Q$  branch for  $\Delta J = 0$  is to be added to the branches of a band corresponding to  $\Delta J = \pm 1$  (positive or  $R$  branch and negative or  $P$  branch). The zero branch can now appear, because the electron transition causes a change of the electric moment, even without a change of the rotational state of the molecule. On the other hand, no change in the electron configuration is connected with a  $\Sigma \rightarrow \Sigma$  transition (page 361); consequently no zero branch can appear in this case, because now the electric moment of the molecule cannot change without a change of the rotational state.

We know from the discussion of the potential curves that the internuclear distance, the moment of inertia, and the binding forces of a molecule may be *greatly* changed by an electron transition. In general, therefore, the sequence of the vibration states as well as that of the rotational states are different in the upper and lower electron states. Accordingly, two different rotational term sequences are shown in Fig. 189, which may be thought of as belonging to any two vibration levels of any two different electron states. The energy difference of the two zero-rotation levels is the wave number of the zero line (forbidden), for which the vibration formula (6-20) applies. It is equal to the sum of the electron and vibration transitions  $\nu' - \nu''$  and is written  $\bar{\nu}_0(\nu', \nu'')$ . The individual rotational lines of a band are designated, as in Fig. 189, by the  $J$  number of the lower state ( $J''$ ). Thus the  $R(3)$  line of a band is the transition  $J' \leftrightarrow J'' = 4 \leftrightarrow 3$ . From Eq. (6-27) we

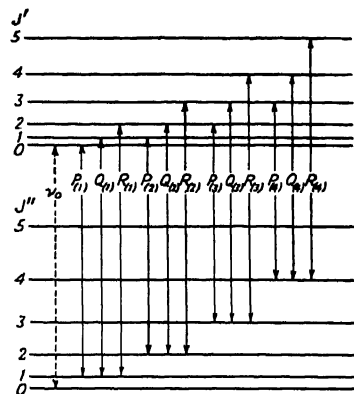


FIG. 189. Transitions between the rotation term sequences of two different electron states of a diatomic molecule. This explains the  $P$ ,  $Q$ , and  $R$  branch normally present in a band connected with an electron transition.

have the formula for the wave numbers of the rotational lines of a band

$$\bar{\nu} = \bar{\nu}_0(v', v'') + B_v J'(J' + 1) - B_{v''} J''(J'' + 1) \quad (6-32)$$

in which, from the selection rule (6-31),  $J''$  must be equal to  $J'$  or  $J' \pm 1$ . These three possible changes in the rotational quantum numbers correspond to the three branches of a band,

$$J' = \begin{cases} J'' - 1 & P \text{ branch} \\ J'' & Q \text{ branch} \\ J'' + 1 & R \text{ branch} \end{cases} \quad (6-33)$$

which are shown in Fig. 189.

The best representation of the structure of an electron band as deter-

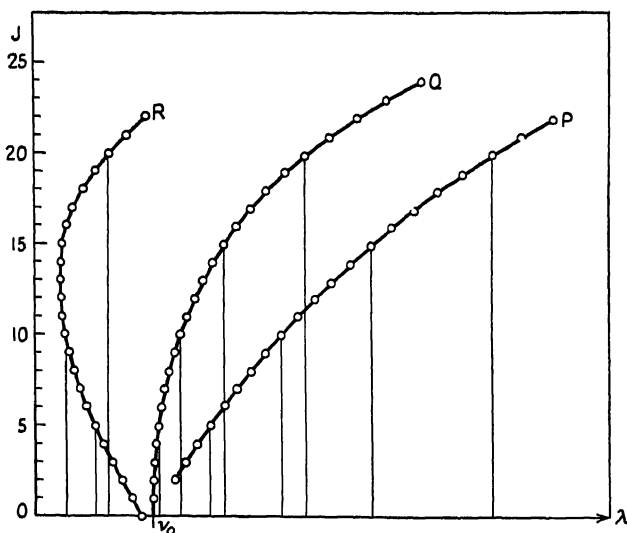


FIG. 190. Fortrat diagram of a band with three branches. Below is the whole spectral band originating from the superposition of the lines of the three branches. (After Fermi.)

mined by (6-32) and (6-33) is its so-called Fortrat diagram. If, as in Fig. 190, the wave numbers of the band lines in reciprocal centimeters or wavelengths are plotted on the abscissa, and the rotational quantum numbers  $J''$  are plotted on the ordinate, the result is a parabola for each of the three branches of the band. This can be easily ascertained by measuring the transition arrows in Fig. 189. The Fortrat diagram thus provides a means of separating the lines of the different branches and rotational quantum numbers which, in the observed spectrum, appear



mixed in a confusing way. Conversely, we can get from the Fortrat diagram a picture of the spectrum (in wave numbers) by projecting on the  $\bar{\nu}$  axis (abscissa) the points in the diagram corresponding to the different lines. This has been done in Fig. 190, where the picture of the whole band is shown underneath the diagram. It is especially clear from the diagram that the band edge is actually not due to a convergence of the band lines, as is the series limit in an atomic spectrum (page 80), but is a more or less accidental phenomenon. It depends, just as the interval between it and the physically significant zero line (for the pure electron

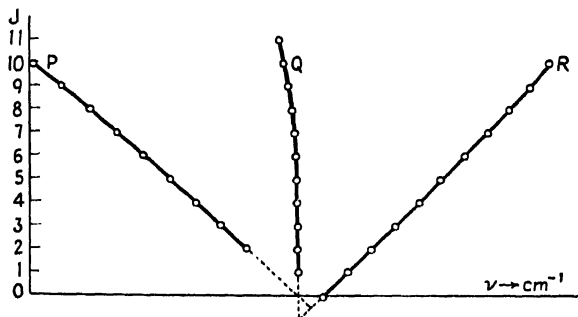


FIG. 191. Fortrat diagram of a band without edge, due to identical internuclear distances in the upper and lower electron state. (After Weizel.)

and vibration transition), on the intervals between the rotation levels of the upper and lower states. If the  $P$  branch forms an edge, as in our Fig. 192, we say that the band is *shaded toward the violet*; in the case of an edge formed by the  $R$  branch (Fig. 190) it is *shaded toward the red*. It follows from the relation between the Fortrat diagram and the transitions in the term diagram that the shading depends on whether the internuclear distance (and with it the moment of inertia) is greater or smaller in the upper state than it is in the lower. We thus have the following relations:

$$\left. \begin{array}{l} \text{Red-shaded band (} R \text{ edge): } r_0' > r_0''; I' > I''; B' < B'' \\ \text{Violet-shaded band (} P \text{ edge): } r_0' < r_0''; I' < I''; B' > B'' \end{array} \right\} (6-34)$$

For equal values of internuclear distances and moments of inertia in the two combining states, we have no edge structure and the  $Q$  branch shrinks together to form a single line as shown approximately in Fig. 191 which is an example of an actually measured band.

Thus a short glance at a band spectrum just to see the shading of the bands allows us to draw important conclusions about the combining molecular states, whereas an exact analysis provides us with the absolute values of the rotation constants, moments of inertia, and internuclear distances.

*d. The Influence of Electron Transitions on the Rotation Structure*

In our treatment of the rotation structure of the electron bands we have not, so far, considered the fact that the total momentum  $\mathbf{J}$  of the molecule which controls the spacing of the rotation terms results from the vectorial addition of the angular momentum of the rotating nuclei and an angular momentum which itself is composed, according to page 361, of the angular momentum of the electron shells about the internuclear axis  $\mathbf{A}$ , and the resultant spin  $\mathbf{S}$ . These relationships become especially complicated by the fact that the rotation of the whole molecule produces a new magnetic field which competes, with respect to the orientation of the spin of the electron shell, with the field in the direction of the internuclear axis. Therefore, the coupling between the angular momenta of the molecular rotation  $\mathbf{K}$ , of the resulting orbital momentum in the direction of the internuclear axis  $\mathbf{A}$ , and of the resultant spin of the electron shells,  $\mathbf{S}$ , changes with increasing rotation, depending on the strength of the internuclear field. The different possible cases of coupling for this interaction of electronic motion and rotation have been analyzed, especially by Hund. These different cases of coupling cause differences in the rotation structure of the electron bands, such as the failure of the appearance of the zero line and of certain neighboring lines, as well as deviations from the parabolic branch curves in the Fortrat diagram. We must pass over all these details of the band structure and can only mention that, as a result of the higher multiplicity of many-electron molecules, multiple branches of bands can also appear. An investigation of all these details, however, is of interest because it permits us to determine empirically the quantum numbers  $\Lambda$  or  $\Omega$  of the electron configuration in the upper and lower states. The complete band analysis thus provides all possible statements about a molecule. The consistent agreement of even the finest details of the often very complicated spectra with the theory is the best proof for the correctness of our theoretical concept of the molecules and their behavior.

*e. The Influence of Nuclear Spin on the Rotation Structure of Symmetrical Molecules. Ortho- and Parahydrogen*

In conclusion we shall briefly consider the influence of the nuclear spin on the rotation structure of diatomic molecules with identical nuclei (homonuclear molecules), upon which an interesting effect, especially in the case of molecular hydrogen, depends. We saw (pages 134 and 233) in discussing the He atom and the  $\text{H}_2$  molecule that, because of the exchange possibility of the two identical electrons, considering

also their electron spin, the term diagrams of these two-electron systems split up into two term systems, a singlet system and triplet system, which do not intercombine. *In the case of  $H_2$  and the other molecules consisting of identical atoms we now have, in addition to the possibility of exchange of electrons, the exchange possibility of the identical nuclei. This again causes, according to the formalism described on page 233, the existence of two term systems which do not combine with each other.* The wave-mechanical treatment showed that the rotation terms (in the electron singlet system as in the electron triplet system of  $H_2$ ) belong alternately to one and to the other of the two term systems which in this case are called even and odd. In the case of  $H_2$ , page 268, we have two nuclear spin moments each of  $\frac{1}{2}(h/2\pi)$ , which can have the same direction (analogous to the triplet system) or opposite directions (analogous to the singlet system). Because of the small interaction between the nuclear spin and the motion of the molecule, the rotation terms which belong to the three possible orientations of the parallel spin moments do not split up. The even and odd terms thus can be distinguished only by their statistical weight, i.e., the rotation band lines corresponding to them are distinguished by their different intensities (in this case 1:3). Actually, as was first found by Meeke, the band spectra of all molecules formed from identical atoms (but, for example, not that of the slightly unsymmetrical  $N^{14}N^{16}$  molecules) show these *alternating intensities of consecutive band lines* which depend on the nuclear spin. In agreement with the theory, the intensity ratio in the case of  $H_2$  is 1:3. This agreement of theory and experiment, and the resulting possibility of determining the spin momentum of the nuclei of molecules from measurements in band spectra, which supplement in a welcome way the nuclear spin determinations by magnetic methods (page 268) and by hyperfine-structure studies in atomic spectra, is an excellent example of the interrelationship of the various fields of atomic physics.

In the case of the hydrogen molecule, whose rotation states are widely separated because of its small moment of inertia, the existence of the two rotation term systems which do not combine with one another leads to an interesting effect. Because this effect is in contradiction to the classical theory, it can be regarded as another proof for the correctness of quantum physics. By continuously decreasing the temperature, not all  $H_2$  molecules can gradually enter the rotationless state  $J = 0$ , as is expected classically. Only the molecules with parallel nuclear spin can enter this state while those with antiparallel spin, because of the Pauli principle, must remain in the next higher,  $J = 1$ , state, which cannot combine with the  $J = 0$  state, regardless of how much we try to cool down the hydrogen. This fact becomes evident in a well-known

anomaly of the specific heat of hydrogen at low temperatures. Thus molecular hydrogen behaves as if it consisted of two different modifications which Bonhoeffer and Harteck called ortho- and parahydrogen. It must be realized, however, that in contrast to the phenomenon of helium, which depends on the effect of the electron spin, the difference between ortho- and parahydrogen depends on the nuclear spin of the atoms which form the molecule, or, more accurately, on the exchange possibility of the two identical nuclei. A transformation of ortho-hydrogen (antiparallel nuclear spin momenta,  $J = 1, 3, 5, \dots$ ) into parahydrogen (parallel nuclear spin momenta,  $J = 0, 2, 4, \dots$ ) would require a change of the direction of one of the nuclear spin moments, which normally is as forbidden for nuclei as it is for electrons (page 147). By special treatment of  $H_2$  under high pressure and by adsorption to cooled charcoal, Bonhoeffer and Harteck and, independently, Eucken and Hiller were able to produce pure parahydrogen. They found, as was expected, that the rotation lines corresponding to  $J = 1, 3, 5, \dots$  were then missing.

#### 6-9. Band Intensities and Temperature Determination by Band Spectroscopy

Now that we have learned about the structure of the spectra of diatomic molecules and its relation to molecular constitution, we shall briefly consider the problem of the intensity of band spectra. This is important, first, because line structure *and* intensity distribution are necessary to give a complete description of a spectrum and, second, because the measurement of the intensity distribution in band spectra has led to the development of an important method for determining high temperatures.

According to page 166, the intensity of a spectrum is determined by the transition probability between the two combining states, and by the number of molecules occupying the initial state of the transition in question, its so-called occupation number. The transition probability, and with it the intensity of the entire band system, is determined by the probability of the electron transitions, computed according to page 216. The transition probability for the individual bands, on the other hand, follows from the Franck-Condon principle (page 370). Within each branch of a band, finally, the transition probability is constant, whereas the relative intensity of the different branches can be computed from  $J$  and the electron quantum numbers according to formulas of Hönl and London. A definite statement about the occupation number of the initial state can be made from Maxwell's distribution law, but only for the case of thermal equilibrium, i.e., in general for the absorption

of molecules. For the case of emission, on the other hand, a statement can be made only if the excitation is due exclusively to the thermal velocity. For the intensity of a spectral line resulting from a transition between states of energies  $E_1$  and  $E_2$  we have

$$I_{E_{12}} = C_{1,2} g_1 e^{-E_1/kT} \quad (6-35)$$

where  $C_{1,2}$  is a constant containing the transition probability between states 1 and 2,  $g_1$  is the statistical weight of state  $E_1$  (see page 166), and  $T$  is the absolute temperature of the absorbing or emitting molecules, i.e., of the initial state.

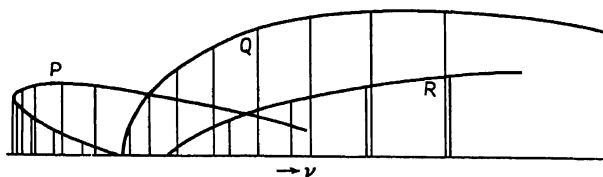


FIG. 192. Intensity distribution in the three branches of a normal band. (After Jevons.)

From the general formula (6-35) we obtain the intensity distribution in a rotation band if we substitute for  $E_1$  the rotation energy

$$E_1 = hcBJ(J + 1) \quad (6-36)$$

and for  $g_1$  the statistical weight of a rotation state

$$g_1 = 2J + 1 \quad (6-37)$$

Equation (6-37) follows from the fact that each rotational state in a diatomic molecule is  $2J$ -fold degenerate as a result of the equal influence of the two rotation axes which are perpendicular to the nuclear axis. With this we have

$$I_J = (2J + 1) e^{-\frac{hcBJ(J+1)}{kT}} \quad (6-38)$$

Equation (6-38) describes the intensities of the rotation lines of a band as a function of the  $J$  values of the initial state. Figure 192 shows the intensity distribution of the rotation lines of a band computed according to (6-38). The intensity of a line is here represented by its length.

According to (6-38) the intensity distribution in a band is temperature-dependent. Thus by measuring the maximum of the intensity in a branch of the band, i.e., the  $J$  value of the line of highest intensity, the temperature of the emitting or absorbing gas can be determined. This method has been applied successfully in measuring the temperatures in electric arcs, especially in the outer zones. However, in discharges

resulting from nonthermal exciting collisions, Eq. (6-38) is not applicable, and attempts to apply it lead to gross errors. Furthermore, by collisions of the second kind with energy exchange (page 104), the occupation of rotation states can be greatly changed, so that all conditions must be carefully checked before band-spectroscopic temperature measurements can be considered reliable.

If the condition of thermal equilibrium is fulfilled, then the same method described above can be used to compute how, at a given temperature, molecules are distributed among the *vibration states* of the initial electronic state. By applying the Franck-Condon principle, page 370, the intensity distribution of the bands within a band system then can be determined. In this case we have to substitute in (6-35)

$$E_1 = hc\omega(v + \frac{1}{2}) \quad (6-39)$$

Since the vibration states of diatomic molecules are not degenerate, the statistical weight of all vibration states is unity. Thus we have

$$I_v = C e^{-\frac{hc\omega(v+\frac{1}{2})}{kT}} \quad (6-40)$$

Because of the difficulty of numerically computing the transition probability contained in  $C$  according to the Franck-Condon principle, Eq. (6-40) is less suited for the *absolute* determination of temperature from the relative intensities of the different bands of a system. However, the determination of *changes* of the temperature from the corresponding changes of the relative intensities of the various bands is easily possible, since in this case  $C$  remains constant. This vibration method has already been used with success in investigating temperatures, especially of electric arcs, and we may be certain that both methods of measuring temperatures by band spectroscopy will be further developed and will play an important role in the development of the physics of high temperatures.

### 6-10. Isotope Measurement by Means of Molecular Spectra

We have already mentioned in the discussion of the phenomenon of isotopes that the detection of isotopes as well as the determination of their masses and relative abundances is possible by optical spectroscopic methods just as by mass-spectroscopic methods. We have treated on page 181 the influence of isotopes on the hyperfine structure of line spectra; now we shall discuss briefly the isotope effect on band spectra.

Since the different isotopes of an atom differ only in their masses, the incorporation of different isotopes in the same molecule (e.g.,  $\text{Li}^6\text{H}$  and  $\text{Li}^7\text{H}$ ) affects the band spectra through the corresponding

change in the moment of inertia (6-29) of the molecule. According to Eqs. (6-27) and (6-28), the separation of the rotational levels is inversely proportional to the moment of inertia  $I$  of the molecule, while according to Eqs. (6-11) and (6-12) the spacing of the vibration levels is inversely proportional to the square root of the moment of inertia  $I$ . The results of these two mass influences on the spectra are the *rotational* and the *vibrational isotope effect*.

We now consider two diatomic molecules, one of which may have the atomic masses  $m_1$  and  $m_2$ ; in the second molecule the mass  $m_1$

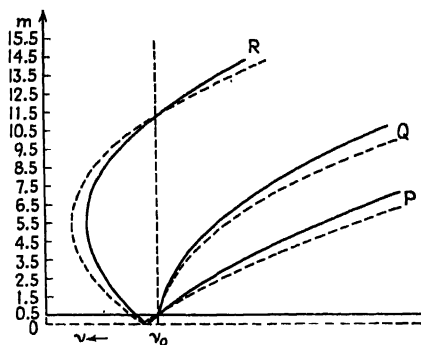


Fig. 193. Rotational isotope effect in an electron band. The branches belonging to the two isotopic molecules are indicated as solid and dotted curves, respectively. (After Mecke.)

is replaced by the mass of an isotope,  $m_1 + \Delta m$ . In computing, by means of the formulas on page 388 for the rotation, the wave-number difference of two equal rotation lines of these two molecules, we find

$$\Delta \bar{\nu}_r = \frac{m_2 \Delta m}{(m_1 + m_2)(m_1 + \Delta m)} \bar{\nu}_r \quad (6-41)$$

where  $\bar{\nu}_r$  is the wave-number distance of the rotation line from the zero line of the band. From the vibration formulas on page 365 we find, for the wave-number difference of two band edges of the two isotopic molecules, the formula

$$\Delta \bar{\nu}_s = \frac{m_2 \Delta m}{2(m_1 + m_2)(m_1 + \Delta m)} \bar{\nu}_s \quad (6-42)$$

in which  $\bar{\nu}_s$  is the distance of the band edge under consideration from the (0,0) band of the band system. Figure 193 shows the rotational isotope effect represented in a Fortrat diagram.

Equations (6-41) and (6-42) show that the isotopic splitting of the lines increases linearly with the distance from the zero line or the (0,0)

band, respectively. By measuring the spacings  $\Delta\bar{\nu}_r$  or  $\Delta\bar{\nu}_s$  and the distances  $\bar{\nu}_r$  or  $\bar{\nu}_s$ , respectively, the mass ratio  $m_1/(m_1 + \Delta m)$  of the isotopic atoms in the molecule can thus be determined. Because of the high accuracy of spectroscopic wavelength measurements this method of determining isotope masses is very accurate. By measuring the intensity ratios of the lines belonging to the isotopic molecules, the ratio of the two kinds of molecules, and thus the relative abundance of the isotopic atoms in the molecule, can be determined very accurately.

By means of this isotope effect in band spectra, new isotopes of the elements C, N, and O were discovered for the first time. A large number of isotopes were confirmed after they had been detected by mass-spectroscopic experiments. Mass ratios and relative abundances have been measured for many new isotopes. Mass-spectroscopic, line-spectroscopic, and band-spectroscopic isotope methods thus supplement each other very nicely.

### 6-11. Survey of the Spectra and Structure of Polyatomic Molecules

The results of molecular research deduced from diatomic molecules as they have been presented in this chapter can be applied, at least in their fundamentals, to polyatomic molecules. However, because of the large number of possibilities for excitation and ionization, for vibration and dissociation, as well as for the rotation about the three axes associated with the different principal moments of inertia, the spectra of these molecules are very complicated. A complete analysis of the electronic, vibrational, and rotational structure, and the corresponding complete knowledge of all molecular data, therefore, is available only for some simpler polyatomic molecules (such as  $\text{H}_2\text{O}$  and  $\text{CO}_2$ ); in most cases we have to be content with incomplete information. Therefore, we shall not discuss in all detail the electronic, vibrational, and rotational structures of polyatomic molecules, but shall limit ourselves to a survey of the more important phenomena which are typical of such molecules.

#### *a. Electron Excitation and Ionization of Polyatomic Molecules*

The systematics of the electron states of polyatomic molecules is closely related to that of diatomic molecules (page 359). It has been advanced especially by Mulliken in the last 10 years. The difference compared to diatomic molecules consists, in the first place, in the larger number of external electrons which can be excited and ionized. In the second place, these electrons, depending on their location in the molecule, play different roles in holding the molecule together. We may distinguish three extreme cases of this behavior.



In many molecules the near ultraviolet absorption is due to electrons which are responsible also for holding the molecule together. For this reason they are called binding electrons. Their excitation often reduces the binding to such an extent (analogous to the case of the diatomic molecule, Fig. 169, page 364) that immediate dissociation into two atom groups results. According to page 377, the corresponding absorption spectrum must then be continuous. This photodissociation, caused by the absorption of radiation by binding electrons, explains why the ultraviolet absorption spectrum of so many polyatomic molecules is continuous. An interjection is necessary at this point. (One must be careful to distinguish, in the case of polyatomic molecules, between genuinely continuous and nongenuine quasi-continuous spectra. The latter are actually discrete band spectra. The reason for this is the following: Since the distance between the band lines according to page 388 is inversely proportional to the moment of inertia of the molecule, in polyatomic molecules with large moments of inertia the distance between the band lines is often so small that the rotation structure cannot be resolved with the usual spectroscopes, even with large gratings. The discrete band spectra then *appear* to be continuous. We cannot discuss here the methods used to distinguish between the quasi-continuous and the genuine continuous spectra.)

In contrast to the first case of the excitation of binding electrons which often produces dissociation, is the case investigated especially by Price, of absorption of light by nonbinding electrons. As in the corresponding case, Fig. 168, page 363 of diatomic molecules, this leads to the second extreme case where the excitation can go over into ionization. In this case the molecule remains stable during the excitation or separation of the nonbinding electron; the nuclear arrangement and vibration are independent of the motion of *this* electron. Spectroscopically, Rydberg series of bands are observed in the vacuum ultraviolet onto which is joined in many cases an ionization continuum, from whose long wavelength limit the ionization energy can be determined. However, it must be remembered that, with polyatomic molecules, one can no longer speak strictly of *one* ionization energy, because the ionization energy varies greatly depending on the binding state of the particular separated electron. Only because, for some reason not yet known to us, *one* electron is preferably excited and ionized, can its ionization energy be called that of the molecule. Our molecular theory here, as so often in atomic physics, has not yet arrived at its final state. We can determine what electron is excited and in what manner, if the molecule absorbs a definite spectrum. However, we cannot yet answer the question why, in a given electron configuration

of a large molecule, one particular electron and only that one electron absorbs the incident light and becomes excited.

A systematic investigation of the third group of excitable electrons, the study of chromophoric groups, can perhaps contribute to the solution of this question. In a polyatomic molecule, radiation can also be absorbed by an electron which does not belong to the shell of the whole molecule but which is strongly localized in a definite group of atoms. This may be a benzene ring or any similar atom group belonging to a large molecular complex. In this case the resulting spectrum is, in a first approximation, independent of the large complex to which the atom group (e.g.,  $\text{NO}_2$  or  $\text{N}=\text{N}$ ) is attached, and only finer detail differences (wavelength displacements) permit one to draw conclusions about the state of binding of the particular group to the whole molecule. Since the absorption spectra of this type of group electrons lie mostly in the visible spectral region and thus determine the color of the material which consists of these molecules, these absorbing groups are called color producers or chromophoric groups.

Our knowledge of the possibility of exciting particular electrons, which depended exclusively on the investigation of absorption, was extended in certain cases by the investigation of fluorescence (page 354). The investigation of the emission spectra of polyatomic molecules for a long time failed to produce reliable results because these complicated structures almost always dissociated in gas discharges as a consequence of exciting electron collisions. What could be observed, therefore, was a confusing superposition of the continuous spectra of the molecule and its components. Recent work of Schüller and Woeldike seems to have opened a new approach for the investigation of the possibility of exciting particular electrons in polyatomic molecules. By using a low-current glow discharge with special precautions, they were able to excite discrete as well as continuous emission spectra of very complicated and easily dissociable polyatomic molecules. By investigating a whole series of such molecules in which either different groups were attached to the same molecule or, conversely, the same groups were attached to different molecular complexes, and also by comparing the emission spectra with the absorption and fluorescence spectra of the same molecule, they were able to establish a large number of empirical rules and relations which seem to be of high importance for understanding the interaction in polyatomic molecules. Interesting differences have been ascertained concerning molecular excitation by light (absorption) and by electron impact (excitation in a discharge). These differences will aid us in the attempt to answer the above-mentioned question why and how a particular electron is preferentially excited in a certain process. In the case of formaldehyde, acetone, and

diacetylene, for example, fluorescent light is emitted from a lower energy state of the molecule than that which is excited by the absorption of light. This initial state of fluorescence is normally reached by collisions of the second kind (page 104) from the state excited by absorption. This same initial upper state of fluorescence which thus *cannot* be reached directly by *absorption*, *can* be excited directly from the ground state by *electron impact*. In other cases, as in the case of benzene and its derivatives, light *absorption* always leads to an excitation of the nonlocalized  $\pi$ -electrons of the benzene ring. Recent investigations have proved, however, that in the glow discharge (i.e., by *electron impact*) this emission of the  $C_6$  ring, which corresponds to a spectrum between 2000 and 3000 Å, is excited only if the mass number of the individual substituents, or that of two substituents coupled by a double bond, is smaller than 27. Since this value is of the order of magnitude of the mass of two C atoms in the ring (mass number 24), we seem to have here a theoretically interesting relation: *the electron excitation can apparently be blocked by externally attached atom groups whose mass is greater than that of the absorbing group*. Similarly, though not quite so pronounced, the mass influence seems to be effective in the case of excitation of the molecules to fluorescence (decreasing intensity of fluorescence in the series  $I_2^-$ ,  $Cl_2$ , Br-benzene). What happens to the energy which has been taken up, if the emission is blocked, seems not yet clear. It does not seem to be transformed to vibration energy. If there exists some unknown storage mechanism, it would be of fundamental interest. With reference to all these questions as well as regarding the detailed behavior of the different outer electrons of larger molecules, many interesting results may be expected from this new experimental approach. In concluding, we mention the interesting observation of Schüller and Woelke that, for example, the sharp emission bands of quinone become diffuse when phenol is mixed with the emitting vapor. Apparently the normal, undisturbed vibration and rotation of the quinone molecule are hindered by the formation of intermolecular "hydrogen bridges," again a very pictorial result.

At the end of this section (page 407) we shall return to the question, which dissociation process is associated with the different observed absorption continua of polyatomic molecules. First, however, we shall briefly discuss molecular rotation.

#### *b. Rotational Structure and Moments of Inertia of Polyatomic Molecules*

Rotational structure, which provided so simple a method for determining moments of inertia and internuclear distances of *diatomic molecules*, is (except for the special case of linear molecules which we

shall discuss soon) of little value in the study of polyatomic molecules. First, the moments of inertia of even the relatively simple polyatomic molecules soon become so large that the rotation bands can no longer be resolved even with highly resolving spectrographs, because, according to page 388, the distance of consecutive band lines is inversely proportional to the corresponding moment of inertia of the molecule. Furthermore, in the most general case, a polyatomic molecule has three principal moments of inertia corresponding to the model of the unsymmetrical gyroscope, and the rotation structure rapidly becomes extremely complicated for this model. Only in the case of the  $\text{H}_2\text{O}$  molecule was

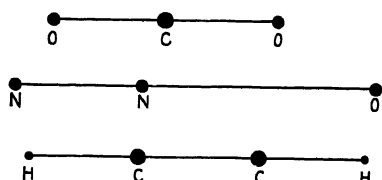


FIG. 194. Schematic representation of some linear polyatomic molecules,  $\text{CO}_2$ ,  $\text{N}_2\text{O}$ ,  $\text{C}_2\text{H}_2$ .

Mecke able to explain almost completely the details of the rotation structure and to determine correspondingly complete molecular data. Even in the apparently simple cases in which two or even three of the moments of inertia are equal, the rotation structure is so complicated because of the interaction between rotation and vibration that only a

few molecules (such as the methyl halides  $\text{CH}_3\text{I}$ , etc., by Herzberg) have been accurately investigated. Actually only the linear molecules such as  $\text{CO}_2$ ,  $\text{N}_2\text{O}$ ,  $\text{C}_2\text{H}_2$ , etc. (Fig 194) have a simple and clearly arranged rotational structure. In these molecules, as in the diatomic molecules, the moment of inertia about the molecular axis is zero in the first approximation (if one neglects the deformation vibration and the resulting orbital momentum of the molecular electrons), and the other two moments of inertia are equal. In this case (Fig. 194) the rotation bands are exactly equal to those of the diatomic molecules so that the formulas of page 388 can be applied. Thus the moment of inertia can easily be deduced from the spectra. However, even if the masses of the atoms forming the molecule are known, in this case the internuclear distances *cannot* be computed without additional information because different nuclear configurations can have the same moment of inertia. Nonspectroscopic methods, e.g., electron or X-ray diffraction, thus must be used to determine the internuclear distances.

### c. Vibration and Dissociation of Polyatomic Molecules

The problem of the vibration of an arbitrary polyatomic molecule formed of  $N$  atoms also implies many theoretical complications. In order to determine the number of possible vibrations we have to subtract from the  $3N$  degrees of freedom of the  $N$  atoms the three degrees

of freedom of translation and in general three degrees of freedom for the rotation of the whole molecule, so that there are  $3N - 6$  degrees of freedom of vibration. Consequently,  $3N - 6$  fundamental vibration frequencies also have to be taken from the spectrum (infrared and Raman). From these, the actual vibrations of the molecule originate (just as Lissajous figures of mechanical vibrations) by superposition with amplitudes corresponding to the different excitation states. These  $3N - 6$  simple (in the first approximation, harmonic) vibrations from which all actual vibrations of the molecule are formed, are called the *normal vibrations*. The six normal vibrations of a plane four-atomic

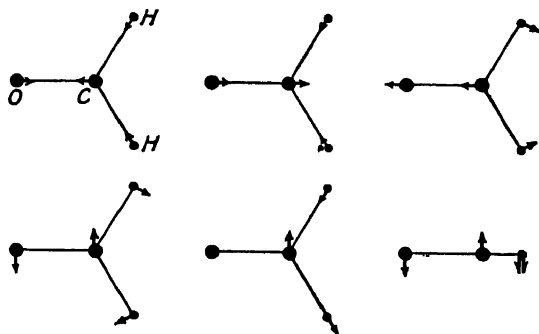


FIG. 195. The normal vibrations of the formaldehyde molecule  $\text{H}_2\text{CO}$ . (After Mecke.)

molecule such as formaldehyde  $\text{H}_2\text{CO}$  are shown, as an example, in Fig. 195. The sixth normal vibration is perpendicular to the plane of the molecule in which the other five vibrations lie.

The fundamental difficulty in investigating the vibration of polyatomic molecules lies in the fact that, in contrast to the diatomic molecules, the spectroscopic determination of the fundamental frequencies of the normal vibrations is *not* sufficient to draw conclusions about the forces between the different atoms and atom groups. One can easily see, for example, from Fig. 195, that the number of spiral and leaf springs one must think of in order to meet the requirements of the normal vibrations shown in Fig. 195, is greater than the six normal vibrations whose frequencies can be determined from the spectrum. The number of force constants is, for  $f$  normal vibrations, equal to  $f(f + 1)/2$ . Consequently, the determination of the normal vibrations from the spectrum is not sufficient to establish a molecular model. In addition, the geometry of the vibrations must be established, and the correct spectroscopically determined fundamental vibration frequencies must be attributed to the different normal vibration forms. The determination of the vibration forms can be made by geometrical considerations in the case of simple

molecules as soon as the nuclear configuration has been established. The vibration isotope effect, page 399, can be used to aid in establishing the form of vibration of complicated molecules. If an isotope of a different mass is substituted for any atom of a molecule, the change in the vibration spectrum produced by this substitution is the greater, the more the substituted atom participates in the vibration. By systematic substitutions the different forms of vibrations and their associated frequencies can thus be determined. Another aid is the fact established by Mecke, that the vibrations along the line joining two nuclei, the "valence vibrations" (e.g., the first normal vibration of Fig. 195), always have a higher frequency than those associated with the angular changes, the "deformation vibrations," such as the fourth and sixth normal vibration shown in Fig. 195. Thus *the difficulty in the vibrational analysis of polyatomic molecules lies not in determining the fundamental vibration frequencies, which are furnished to a large extent by Raman spectroscopy, but in attributing these frequencies to the normal vibrations of the vibration forms.*

It is self-evident that the vibrations of polyatomic molecules are anharmonic, and thus our treatment of page 367f. is valid also for the individual normal vibrations. In particular, each valence vibration can be represented by a potential curve such as Fig. 172, and this representation is of value in explaining the dissociation processes which we shall discuss below. From the anharmonicity of the vibration it follows, as in the case of diatomic molecules, that, in addition to the transitions between neighboring states, larger vibrational transitions also occur in the spectrum. However, these appear with smaller intensity. Moreover, it follows from the vibration anharmonicity that there is a mutual influence of the different vibrations which is manifested in the appearance of *combination vibrations*, i.e., the superposition of different normal vibrations. Obviously, this leads to an increased complication of the vibration spectrum. This situation has an analogue in acoustics where the analysis of a composite sound phenomenon, as from an organ pipe, offers many difficulties. Just as in the spectra of diatomic molecules, the frequencies of all those vibrations which cause a change of the electric moment of the molecule occur in the emission and absorption spectra of polyatomic molecules. The frequencies of those vibrations which change the polarizability of the molecule can be taken from the Raman spectrum. Again infrared and Raman spectra supplement each other in supplying the desired information, so that almost all molecular vibrations can be determined experimentally.

We distinguished above, in discussing the motion of electrons in polyatomic molecules, between localized electrons associated with a

definite atom or atom group, and nonlocalized electrons associated with the shell of the whole molecule. In an analogous way, we can distinguish between *localized* and *nonlocalized vibrations* according to whether the energy of the vibration quantum is dependent upon the nature and the state of motion of its environment in the molecule or not. For example, according to Meeke, the vibration of two doubly bound C atoms in all hydrocarbons has the same fundamental vibration quantum  $\omega_0$  of about  $1,600\text{ cm}^{-1}$  and so it can be termed strongly localized, whereas the *C—H* vibration varies from  $2,750$  to about  $3,150\text{ cm}^{-1}$  with increasing binding value of the C atom and consequently is somewhat dependent upon the environment, thus it is not localized. Such empirical rules, determined from investigations of a series of homologous molecules, are of great value for the understanding of polyatomic molecules.

Of much greater importance for polyatomic than for diatomic molecules is the interaction between vibration and rotation. It is evident from the vibration forms such as those in Fig. 195 that this vibration will be more or less distorted by molecular rotation, depending on the particular type of excited vibration. This is analogous to the perturbation of the vibration of a Foucault pendulum by the Coriolis force due to the earth's rotation. Finally there are certain cases when a molecular vibration can go directly over into a rotation. This is possible if, for example, a molecular group such as  $\text{CH}_3$  is either free to rotate with respect to the rest of the molecule or is held in its stable position by such weak forces that it can describe slow rotatory vibrations about its rest position. Occasionally these can go over into a real rotation. Thus the clear distinction between vibration and rotation, which is characteristic for the case of diatomic molecules, can get completely lost in the case of some polyatomic molecules. We cannot discuss here the correspondingly complicated spectral phenomena.

After this survey of the vibrations in polyatomic molecules we finish our discussion by briefly considering the question of dissociation processes. As in the case of diatomic molecules, dissociation by absorption of continuous radiation due to excessive vibration is possible (page 377), but it is limited naturally to the overexcitation of valence vibrations (breaking of bonds). This clearly emphasizes the significance of valence vibrations as compared to deformative vibrations. *Because of the large number of possible vibrations, there is a large number of possibilities for the molecule to dissociate.* For this same reason, there appears much more frequently than in the case of diatomic molecules a superposition of discrete and continuous vibration term systems which, according to page 388, can lead to predissociation. In both cases,

consequently, the large number of possibilities makes it often very difficult to attribute a particular absorption continuum or a wavelength of predissociation in the spectrum of a polyatomic molecule to a definite dissociation process.

In favorable cases, photochemical data can be used to compute approximate dissociation energies for the different possibilities of dissociation, and then to attribute them to the observed continua (or their maxima). This was first done by Schumacher, Stieger, and the author when they associated the longwave maxima of the pure continuous  $\text{Cl}_2\text{O}$  absorption spectrum with its decomposition into  $\text{ClO} + \text{Cl}$  with and without excitation of the dissociation products, whereas they attributed the shortwave maximum to a photodissociation of the molecule into the three atoms  $\text{Cl} + \text{Cl} + \text{O}$ . This interpretation was confirmed by the simultaneous investigation of the photochemical decomposition. Upon irradiation by light of the different wavelengths, a different quantum yield was found which agreed with that expected for the different processes. Similar investigations also seem promising for other molecules. In certain cases, furthermore, a direct chemical detection of dissociation products may be possible and could be an aid in checking the correctness of an assumed decomposition process. Occasionally one can go farther, as was first done by Wieland in interpreting the continuous absorption spectrum of the metallic halides of the type  $\text{HgI}_2$ . He was able to identify the excited dissociation product by its emission and thus confirm the particular dissociation process. Finally, by investigating the absorption spectra of an entire series of homologous molecules and by substituting other atoms at various places in the molecule, the decomposition processes can be clarified and associated with their continuous spectra or regions of predissociation.

In general we have to understand that, with polyatomic molecules, the value of the dissociation energy determined from the long wavelength limit of an absorption continuum (or the wavelength where predissociation begins) indicates only an upper limit. Because of the coupling between the molecular vibrations, frequently, in the process of dissociation by the absorption of light, some of the energy is used to excite other vibrations.

To summarize: The investigation of the spectra of polyatomic molecules will give us answers about the excitation, ionization, and dissociation processes and the amounts of energy required for these processes. In the simpler cases it also permits us to determine the possibilities of vibration and the force constants as well as the moments of inertia. For all complicated molecules, however, the determination of the nuclear configuration and the binding forces is possible only if the spectroscopic



methods are supplemented by the other methods of molecular physics mentioned on page 348, especially by electron and X-ray diffraction.

### 6-12. The Physical Explanation of Chemical Binding

In this chapter we have discussed the most important facts about the structure and properties of molecules. In doing this, we did not consider the difficult problem, which atomic forces account for the binding of a number of atoms, either like or unlike, in a molecule. This is really the fundamental problem of chemistry, because otherwise chemistry has to accept simply as a sort of miracle the fact that there is an  $H_2$ , but no  $H_3$ , molecule, a  $CO$  and a  $CO_2$ , but no  $CO_3$ , an  $H_2SO_4$ , but no  $HSO$ , etc. The theoretical explanation of the periodic table (page 168) was the first major achievement of atomic physics in basic chemistry. The physical explanation of chemical binding, achieved by the methods of quantum mechanics in 1927, completed the successful attempt to understand the basic features of chemistry from the properties of the atoms.

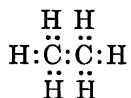
Partial success had been obtained, as early as 1916, by Kossel, who succeeded in explaining the binding in heteropolar or ionic molecules (page 381), such as  $NaCl$ , as an effect of electrostatic forces. The Bohr theory, just developed at that time, explained the unique behavior of the chemically inactive noble gases from their completed (saturated) electron shells. Atoms with one or a few external electrons outside of closed shells (such as alkali or alkaline-earth atoms) are electropositive, because they can attain noble-gas configuration by release of these external electrons. The halogen atoms, on the other hand, are electronegative, because they can complete their noble-gas shell by receiving an additional electron. The binding of an alkali-halide molecule, according to Kossel, is based on the following: A  $Na$  atom, for instance, gives up its external electron to a  $Cl$  atom, thus producing a  $Na^+$  and a  $Cl^-$  ion. These ions, which are very stable because of their ideal noble-gas configuration, attract each other electrostatically and thus form a polar molecule.

Also polyatomic molecules can be explained by this effect of ionic binding. Of the combining atoms, those from the left side of the periodic table always appear positively charged, those from the right side negatively charged. Atoms from the central groups of the table can appear as positive or negative ions depending on whether they combine with atoms which are farther to the right or farther to the left in the periodic table. The molecules  $PH_3$  and  $PCl_5$  are examples of this. In  $PH_3$  the electron shell of the  $P$  atom is completed, by inclusion of the three hydrogen electrons, to form the argon noble-gas shell. It then

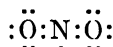
becomes a threefold negatively charged  $P^{---}$  ion which binds the three singly positive  $H^+$  ions. In the  $PCl_5$  molecule, on the other hand, each of the five Cl atoms tears away an electron from the  $P$  atom, so that its electron shell is reduced to the noble-gas shell of neon and we get a fivefold positive  $P^{+5}$  ion which binds electrostatically the five  $Cl^-$  ions. In an analogous way, we may regard the  $C$  atom in  $CH_4$  as a fourfold negatively charged, and in  $CCl_4$  as a fourfold positively charged ion. Abegg's rule (1906), according to which the sum of the highest positive and negative valency of an atom is always 8, thus follows automatically from the possibility of completing or reducing the electron shell to that of a noble gas with eight outer electrons.

Kossel's explanation of polar binding is clearly limited to the binding of atoms whose external electron shells supplement each other to form noble-gas shells. The binding of the simplest existing molecules, such as  $H_2$ ,  $N_2$ ,  $O_2$ , and many others, i.e., of the so-called homopolar molecules, cannot be explained on this basis.

An attempt at an explanation was made somewhat later by G. N. Lewis with his octet theory which is still very popular with many chemists. Lewis's theory is also based on the fact that all electron shells tend to complete the noble-gas shell of eight electrons, so-called octets. By indicating each external electron (i.e., electrons outside of closed shells) by a dot, Lewis writes the ethane molecule  $C_2H_6$  in the following way:



By making use of the H electrons, the C atoms thus succeed in surrounding themselves by complete octets. Each electron pair, however, belongs to *two* atoms and thus is supposed to cause the binding between them. For many well-known molecules the octets cannot be complete.  $NO_2$ , for instance, must be written as



However, Lewis was able to show that such molecules are the less saturated, i.e., the more active chemically, the less complete their octet shells are. It is an important feature of Lewis's concept that he recognized homopolar chemical binding as an effect of electron pairs common to the two bound atoms. He could not yet give, however, an *explanation* for this binding effect of the electron pairs.

This explanation was given, on a quantum-mechanical basis, in 1927 by Heitler and London for the  $H_2$  molecule. According to them,

homopolar binding is caused by the exchange forces between the two H atoms; the nature of these forces has been treated on page 228 for this very example. If the distance between two H atoms, assumed to be far away from each other in the beginning, is gradually diminished, and the interaction of the electron clouds with the two nuclei is taken into account, a shallow minimum of the potential energy of the system, i.e., some binding of the atoms, results from purely classical electrostatic forces. However, the binding energy accounts for only a small fraction of the actually known binding energy of the  $H_2$  molecule. We learned on page 234 that the quantum-mechanical binding is a consequence of an exchange resonance. In the unperturbed state (large distance between the two H atoms) the system is described by two eigenfunctions which belong to the same energy eigenvalue, because the two electrons are completely identical. This system is perturbed if the distance between the atoms is reduced until the electron eigenfunctions begin to overlap. The perturbation theory for this case shows that the two energy states of the same energy split up just as do the levels of the He atom in singlet and triplet terms. As a consequence of this interaction, a repulsive state of the system of two atoms is formed whose energy increases steadily with decreasing internuclear distance, and a "bound state" whose energy decreases to a minimum and then increases again with decreasing internuclear distance (Fig. 119). This potential minimum of the singlet state corresponds to the bound  $H_2$  molecule. The numerical calculation of the binding energy and the internuclear distance of  $H_2$  by Heitler and London did not agree too well with the data derived from the band spectra. Further refinements of the computations, however, have resulted in complete agreement with the empirical data. We have seen on page 232 that the eigenfunction of the bound state in a first approximation is obtained by *addition* of the eigenfunctions of the unperturbed system of the two atoms, whereas that of the repulsive state is obtained by subtraction [Eqs. (4-123) and (4-124)]. Explained pictorially, this result means that binding occurs between two atoms if the electron density between the nuclei (given by the norm of the eigenfunction) is large. Moreover, the bond is the stronger, the more the electron eigenfunctions overlap, because the value of the exchange integral (4-121) depends on this overlapping. We shall make use of this important result in the very pictorial explanation of the angular valences of stereochemistry.

We have seen on page 233 that, as a consequence of the Pauli exclusion principle, the spin directions of the two electrons in the stable ground state of the  $H_2$  molecule must be antiparallel, whereas they are parallel for the repulsive state. For this reason, the magnetic saturation

of the antiparallel spins is occasionally described as the essential feature of homopolar binding. The name *spin valence* is used in this connection. However, *not the spin saturation, but the sign and value of the exchange integral* [(4-121) of page 231] *are decisive for the binding*. In general, we have antiparallel spin directions in the binding case, because that leads to an electron eigenfunction which is symmetrical in the electrons, so that for antiparallel spin directions we find a maximum of the electron density between the nuclei. That spin compensation is not a decisive effect may be seen from the case of the  $O_2$  molecule. Its  $^3\Sigma$  ground state is evidence that the spin directions of the two outermost electrons are parallel. In this case the *orbital* momenta  $L = 1$  of the oxygen atoms (ground state  $^3P$ ), instead of the spin momenta, compensate to form the  $^3\Sigma$  ground state of the  $O_2$  molecule. This is called *orbital valence* by Heitler.

Exchange forces, which here are considered responsible for homopolar binding, are known to us from the binding of protons and neutrons in the nucleus (page 337). Just as nuclear binding could be regarded as a consequence either of an exchange of protons and neutrons, or as a consequence of a (virtual) emission and absorption of mesons by the nuclei (page 337), we can regard the exchange forces of the molecular bond as due to an exchange of electrons, or as caused by (virtual) emission and absorption of photons.

Because of the basic difficulty of a pictorial concept of the exchange forces, it is simpler to approach the whole problem of chemical binding from the energy point of view. Chemical binding between two atoms occurs and can occur only if the system of the two atoms has, at a small internuclear distance, a state of lower potential energy than if the atoms are completely separated. Herzberg, Hund, and Mulliken have developed this approach. They consider the energy of the molecule under consideration as composed of the contributions of all individual electrons, which are thought of as being built into the fixed arrangement of the nuclei of the molecule. The problem is now whether energy has to be added or is released when a particular electron is transferred from the distant atom into the molecule, i.e., whether this electron is a "binding" or a "loosening" electron. This study is confined to only the valence electrons of the atoms which form the molecule, because the internal electron shells are saturated and therefore do not participate in the chemical bonds. A certain particular bond then can be characterized simply by the difference of the number of binding and loosening electrons. In the  $O_2$  molecule, for instance, three of the four valence electrons of the two O atoms find places in binding electron states of low potential energy ( $2p\pi^4$  and  $3p\sigma^2$ ), whereas the fourth electron pair has to use up

energy in order to reach the only available higher state ( $3p\pi^2$ ). This electron pair thus has a loosening effect and compensates the binding effect of one of the binding electron pairs. The final result thus is that we have two uncompensated binding electron pairs which are responsible for the double bond known to the chemist for the  $O_2$  molecule.

In an analogous way it is easy to find that two atoms with two  $s$  valence electrons (e.g., the atoms of the second group of the periodic table) cannot form a molecule, because only one of the electron pairs can go into a lower energy state of the molecule, whereas for the second pair only a higher energy state is available. The effects of the binding and the loosening electron pair thus compensate each other, and a strongly bound homopolar  $Hg_2$  molecule cannot exist, in contrast to the loosely bound  $Hg_2$  van der Waals molecule, which we discussed on page 385. It is obvious that this method of investigating the binding requires an exact knowledge of the molecular electron states, whose importance thus is nicely demonstrated to the beginner. It allows us, of course, to study theoretically all possibilities of binding between given atoms, i.e., makes possible an investigation of chemically unknown molecules.

Finer details of the possible bonds and their characteristics can be discussed if we make use of the wave-mechanical "shape" of the valence electrons, i.e., their eigenfunctions according to Fig. 116 (page 227). The  $p$  electrons here play a particularly important role because, in contrast to the  $s$  electrons with their spherical symmetry, they have an axial symmetry. The three  $p$  electrons of the nitrogen atom, for instance, are oriented in such a way that their axes are mutually perpendicular. In the  $H_2O$  molecule, for example, the two  $s$  electrons of the H atoms are bound by two  $p$  electrons of the O atom (bond type  $s-p^2-s$ ). As these two  $p$  electrons are oriented perpendicularly to each other, it is evident that  $H_2O$  cannot be a linear molecule. According to Fig. 196, the  $p$  electron along the  $x$  axis can bind only an H atom whose nucleus is also on the  $x$  axis, whereas the second  $p$  electron of the O atom, because of its orientation along the  $y$  axis, can bind only an H atom whose nucleus is near the  $y$  axis, for only then a sufficient overlapping of the electron eigenfunctions is possible. *From the wave-mechanical "shape" of the  $p$  electrons, i.e., from the  $|\psi|^2$  distribution in space, we have thus deduced the important result that the  $H_2O$  molecule cannot be a linear molecule but should have a right angle.* In a similar way, we get, quite generally, the angular valences of stereochemistry.

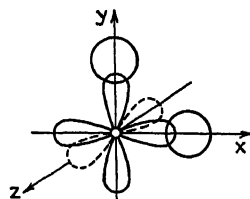


FIG. 196. Schematic representation of the binding in the  $H_2O$  molecule by  $p$  valences. (After Stuart.)

The band analysis of  $\text{H}_2\text{O}$ , carried out first by Mecke, has led to an angle of somewhat over  $100^\circ$ , and similar deviations from right angles are found quite generally in stereochemistry. These deviations can be explained in a very pictorial way by the assumption that the atoms at the sides of the angle interact with each other and thus increase or decrease the valence angle.

An example of the binding of three  $s$  electrons by the three mutually perpendicular  $p$  electrons of another atom is the ammonia molecule,  $\text{NH}_3$ . Its pyramidal shape thus is easily understood to be a result of the binding of one H atom on each of the  $x$ ,  $y$ , and  $z$  axes, and an additional attraction among the H atoms.

More complicated is the case of the tetravalent carbon atom which is of decisive importance for all organic chemistry. We shall discuss briefly the simplest organic molecule, methane,  $\text{CH}_4$ . According to Table 10 (page 178), the C atom has, outside its closed  $K$  shell, two  $2s$  electrons and two  $2p$  electrons. It happens, however, that these four valence electrons of the carbon atom have approximately the same energy. We speak of an approximate  $s$ - $p$  degeneracy. By their mutual interaction, the difference between  $s$  and  $p$  electrons disappears to such an extent that the C atom has four practically equivalent valence electrons. As a consequence of general symmetry rules, the  $\text{CH}_4$  molecule then should be, in agreement with experience, a regular tetrahedron with the C atom in its center. The particular role which the C atom plays in organic chemistry thus seems to be due to a peculiarity of its electron configuration, the approximate degeneracy of the two  $s$ - and the two  $p$ -valence electrons. On page 410 we have treated the same  $\text{CH}_4$  molecule on the basis of the Kossel theory of polar binding. We therefore should like to use this opportunity to point out that polar and homopolar binding are not nearly such fundamentally different effects as it might appear at first sight. If, in the time average, one of the two exchanging electrons is always near the atom  $A$ , and one near the atom  $B$ , we have a regular atomic molecule and so-called homopolar binding. If, on the other hand, both electrons which are responsible for the bond are predominantly near the atom  $A$  and none near the atom  $B$ ,  $A$  appears to be negatively and  $B$  positively charged.  $AB$  is then a polar molecule (such as  $\text{NaCl}$ ), and we speak of polar or ionic bonds. Thus polar binding also can be understood on the basis of the exchange concept. Between these two limiting cases there exist innumerable transitions, i.e., molecules  $AB$  in which the two binding electrons are not always, but more often, near atom  $A$  than near  $B$ , so that we might speak of partial ionic and partial homopolar binding. We shall not go into any detail; we want to indicate only the fact that basically every

bond can be understood by one theory, and that ionic and homopolar bonds are only the two limiting cases.

The different bonds discussed so far are designated as *localized bonds*, because every electron pair can clearly and definitely be attributed to a bond between two atoms and thus causes the bond between them. This concept of localized binding seemed to be the only possible one. However, the most important molecule of organic chemistry, the benzene ring  $C_6H_6$ , could not be explained on this basis. If one of the four valence electrons of each C atom binds one H atom, each C atom has three free valence electrons. In the benzene ring we thus should find single and double bonds alternating between the C atoms, and this actually was the opinion of its discoverer, Kekulé. This hypothesis, however, is in contradiction to the chemical and physicochemical evidence which points clearly to a completely symmetrical ring. E. Hückel has solved this problem by wave-mechanical methods. He was able to prove that each C atom is bound to its two neighboring C atoms by single bonds, and that *the remaining six valence electrons of the six C atoms are distributed uniformly over the whole benzene ring*. We thus might speak of one and one half bonds between each two C atoms. As there are no one and one half electron pairs, it is more correct to speak of one electron pair (one bond) between every two C atoms, whereas the remaining six electrons are called "*nonlocalized*" *valence electrons*. In our pictorial concept of electron exchange it seems quite reasonable to assume that electrons can not only be exchanged between two neighboring atoms (localized bond), but that the six nonlocalized electrons exchange their places in such a way that in the time average always one electron is near each of the six C atoms of the benzene ring. Quite a number of additional details of chemical bonds in organic molecules can also be explained by quantum mechanics, among them the rigidity of the double bond against torsion.

The atomic valence theory in its quantum-mechanical form thus is able to account for all details of the fairly complicated phenomena of chemical binding. Statements can be made also about the binding character of unknown molecules. No doubt, the mathematical difficulties of the perturbation theory as applied to large and complicated molecules become so insurmountable that in practical cases we have to confine ourselves to qualitative studies and statements. Nevertheless, the basic solution of the problem of the chemical bond is of highest importance for theoretical as well as practical chemistry. For this reason it may be especially well suited to demonstrate to the chemist the necessity of familiarizing himself with atomic physics, including the systematics of electron states and their quantum-mechanical treatment.

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## CHAPTER 7

### ATOMIC PHYSICS OF THE LIQUID AND SOLID STATE

#### 7-1. General Review of the Structure of the Liquid and Solid States of Matter

We now turn from the polyatomic molecules to the consideration of the very large atomic and molecular complexes, i.e., to liquids and solid bodies, though the liquid state can be treated only very briefly in this section. Because of the symmetry of the binding forces between atoms in the equilibrium state, i.e., in the minimum of potential energy, the atoms or molecules arrange themselves in a very symmetrical geometric pattern or lattice which we call a crystal. For weaker binding forces or at higher temperatures, the thermal motion of the atoms prevents them from attaining this state of equilibrium. The atoms or molecules then have a definite mobility and the regularity of the atomic pattern is more or less disturbed. This is the picture of the liquid state of matter.

*The solid state, consequently, in its ideal form is identical with the solid crystalline state.* In the apparently noncrystalline or amorphous structures we have either a large number of coalesced microcrystals or, as in the case of glass, an undercooled liquid which is not a genuinely solid body at all. The very few solids which are regarded as "really amorphous" at the present time, such as the "explosive" modification of antimony and one modification of selenium, apparently are not *stable* modifications of the solid state, as is evidenced by the explosibility of the amorphous antimony. The rest of this chapter will be devoted to a thorough discussion of the solid crystalline state. In this section, therefore, we shall take up briefly the problem of the structure of the liquid state, of which not too much is known at present.

From the above interpretation it is clear that in liquids, in which the average interatomic distances, because of their equivalent densities, are essentially the same as in crystals, the forces between atoms or molecules also tend to establish a crystallike, regular arrangement of the atoms. However, the crystalline structure is disturbed more or less by the thermal motion of the atoms, molecules, or ions. This explains that in liquids we find phenomena of order as well as of disorder. For the same reason, a completely random distribution of atoms cannot

even be expected in compressed gases, and evidence of some degree of order in such gases has been found by the X-ray method which we shall discuss below. Although for a long time it had been assumed that the lack of a geometrically regular molecular arrangement was a characteristic of liquids, recently the indications of order have been studied in more detail. It has been concluded from this evidence that there is good reason to speak of a semicrystalline state of fluids, as O. Lehmann did so vividly 50 years ago in his "World of the Liquid Crystals."

*What at first appears to be a radical distinction between the solid crystalline and the liquid state, is evidently a difference in the degree of mobility of the constituents in the crystal lattice which more or less retains its regular structure.* We shall see on page 454 in the discussion of diffusion in crystals that even in the most ideal solid crystals the constituents have a certain mobility, and that a number of very important effects depend upon this mobility. This mobility of single lattice constituents (and with it the general disorder of the crystal) increases with increasing temperature until at a definite temperature the melting point is suddenly reached, above which a certain mobility is a characteristic property of all particles. Even in the fluid state the mobility increases more and more with increasing temperature, and this must cause a gradual disappearance of the crystalline properties. That we do not have a gradual transition from the solid to the liquid state, but a sharp melting point, is due to the fact that only the solid and the liquid state are thermodynamically stable, but not the intermediate states, where the solid crystal begins to loosen up appreciably.

Because the assertion of the semicrystalline structure of liquids is so surprising at first, we shall briefly discuss the most important evidence in favor of it.

Debye's method of scattering of X-rays by liquids provides one direct proof. If the distribution of molecules and the intermolecular spacing were completely random, the scattered intensity would decrease uniformly with increasing scattering angle. However, if there is a regular molecular arrangement, then maxima and minima are to be expected as a result of the interference of the X-rays scattered by the regularly spaced molecules. This was actually observed. If for a number of possible geometrical arrangements (i.e., liquid structures) the theoretically expected scattering curves are deduced and these compared with those obtained empirically, the molecular structure in a liquid can be determined with considerable accuracy in those cases which are not too complicated. Recently the atomic arrangement in liquid mercury has been studied in this way by Hendus who used monochromatic X-rays. At 18°C the atomic arrangement agreed essentially with that of crystal-

lized mercury, and thus deviated considerably from the closest spherical packing which was expected formerly for the liquid state of mercury.

A no less significant evidence of the semicrystalline structure of liquids is the fact that the measured value of the atomic heat of monatomic liquids such as mercury or liquefied argon is 6 cal/mole. This is twice the value to be expected for freely mobile atoms each of whose three translational degrees of freedom contributes  $R/2 = 1$  cal/mole to the atomic heat. On the other hand, the constituents of a crystal vibrate about a rest position, so that  $3R/2$  cal are contributed to the atomic heat for the kinetic vibration energy and an equal amount for the potential energy, since in a harmonic vibration half of the total energy is on the average kinetic and half potential energy. Consequently the atomic heat of a solid body (if all degrees of freedom are excited) is  $2(3R/2) = 6$  cal/mole. That this same value was found for monatomic liquids can only be explained by the hypothesis that their atoms vibrate in three dimensions about a rest position, except that in the case of liquids this center of vibration itself participates in a random translatory motion which is temperature dependent.

There are a number of further results of optical and electrical measurements, which we cannot discuss here, that indicate just as clearly the crystallike arrangement of the molecules in liquids. Stuart and Kast have investigated the dependence of liquid structures on the type of molecule, on the temperature, and density by model experiments. Their model liquids consisted of large numbers of differently shaped little magnets, which were shaken continuously as a substitute for the thermal motion, and in which the intermolecular forces were replaced by the forces between the small magnets which replaced the molecules. The different structures of dipole- and quadrupole-molecule liquids and quite a number of finer details can be shown very nicely in this manner.

Generally the existence of strong dipole or quadrupole moments introduces a complication in the normal liquid structure due to the formation of molecular chains (e.g., in alcohols) or molecular clusters as in the case of the so-called associated liquids whose anomalous behavior depends on the formation of these clusters. By far the most important associated liquid is water whose anomalous behavior has long been attributed to association. It had been thought earlier that water consisted of polymeric molecules of the type  $(H_2O)_n$  with a constant degree of association  $n$ . However, all attempts to determine  $n$  have failed. From this failure and other evidence, the conviction has grown that water consists of molecular clusters of indefinite size. The exact structure of these molecular clusters can be determined by comparing the

experimental X-ray scattering curves with the theoretically deduced patterns. By this method a tridymite-like structure was found in which, according to Fig. 197, each O atom is tetrahedrally surrounded by four protons. This particular geometric arrangement of the  $\text{H}_2\text{O}$  molecules, *i.e.*, this *semicrystalline structure of water*, is responsible for the well-known abnormal properties of water. From this, the change in the structure and properties of water as a result of the addition of relatively few ions (or the addition of a little alcohol) can also be understood: the tridymite structure of water which depends on secondary binding forces (van der Waals forces, page 386) is greatly distorted by the electrostatic forces of the ions or by the addition of a few large foreign molecules, which at the very least influence the size of the molecular clusters. Conversely, the addition of a few  $\text{H}_2\text{O}$  molecules to pure alcohol cannot noticeably alter its chain structure. In agreement with this, hardly any change in the properties of alcohol is noticeable when it is mixed with small amounts of water. Many empirically known properties of liquids can be understood in this way from their crystal structure, *i.e.*, from the viewpoint of atomic physics.

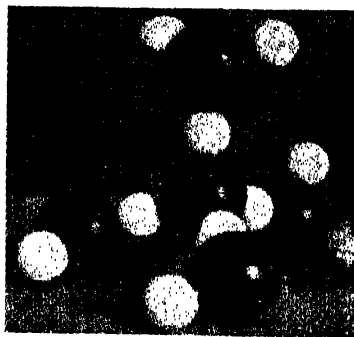


FIG. 197. The quasi-crystalline tridymite structure existing in liquid water due to association of the  $\text{H}_2\text{O}$  molecules. Black spheres = oxygen atoms, white spheres = hydrogen atoms. Each O atom is tetrahedrally surrounded by four H atoms. (After Kortüm.)

So far as can be concluded from the very incomplete investigations, in general, the quasi-crystalline structure of liquids is the more pronounced the greater the dipole moment or the larger and more complicated the individual molecules. Evidence for this can be derived from many remarkable phenomena which Lehmann found in his studies of large and complicated organic molecules, described in his book on liquid crystals. Most surprising among his results is the double refraction of certain liquid droplets and, in many cases, a very pronounced dichroism (different absorption in different directions of orientation). Both of these phenomena are clear evidence for an anisotropy which is characteristic of a crystal. On the other hand, the internal mobility, *i.e.*, the change of the shape of the drops, was just as evident. Therefore, Lehmann's designation as "liquid crystals" seems correct in spite of the objection of many crystallographers. Thus we have to distinguish between the special phenomenon of liquid crystals and the general phenomenon of the solid crystals. A discussion of Lehmann's results from

the standpoint of modern atomic theory, unfortunately, is still lacking in spite of the recently increasing interest in the structure of liquids.

With these remarks we close our brief discussion of the liquid state. Whenever in the following the term crystal is used, it applies exclusively to the normal, solid crystals.

## 7-2. Ideal and Real Crystals. Structure-sensitive and Structure-insensitive Properties of Solids

Physics of the solid state is more or less identical with crystal physics. We begin our discussion by stressing the important difference between *ideal* and *real* crystals. The ideal rock salt crystal, for example, consists of negative chlorine ions and positive sodium ions arranged in a regular geometric lattice in which each sodium ion is surrounded by six chlorine ions, and conversely. However, this ideal crystal is only a *model* and does not actually exist, first, because crystal growth does not occur quite regularly, but occasionally there is a lattice defect (e.g., a missing constituent ion) and, secondly, because the purity of the crystal material is not so high but that here and there an impurity atom is incorporated in the crystal lattice. Actually, even the best real crystals have at least one missing ion or impurity atom among  $10^6$  normal constituents, i.e., according to page 18, at least  $10^{16}$  lattice defects per cubic centimeter.

Smekal first pointed to the fact that an important difference in the properties of crystals is related to this distinction between ideal and real crystals. He designated those properties of crystals which were essentially independent of lattice defects or impurity atoms as *structure-insensitive*. Among these structure-insensitive properties we may name the lattice structure, the specific heat, elasticity, thermal expansion and compressibility, the energy of formation, the principal features of optical absorption (color) and dispersion, as well as, for metallic crystals, the electronic conductivity, and finally dia- and paramagnetism. Those properties of crystals, on the other hand, which depend essentially on the number, arrangement, and type of lattice defects and impurities, are called *structure-sensitive*. Among these are especially the diffusion phenomena, the ionic and electronic conductivity in insulating crystals and semiconductors, the internal photoeffect, plasticity, and crystal strength, as well as certain finer features of optical absorption and luminescence which are of importance for the typical structure-sensitive phenomenon of phosphorescence.

The crystal defects consist, on the one hand, of lattice flaws, most of which are either holes in the lattice or lattice atoms outside their regular geometrical position (interstitial atoms), and, on the other hand,

of impurity atoms in normal or abnormal lattice points. A real crystal occasionally consists of many small, almost ideal microcrystals of about  $1\mu$  diameter with defects near and on the boundaries between the individual microcrystals (layer or mosaic structure of a crystal).

In the following we shall treat first the ideal crystal and the structure-insensitive properties of solid matter. In the second part of this chapter we shall then discuss the real crystals and those phenomena and properties of solids which are caused by lattice defects and impurities.

### 7-3. The Crystal as a Macromolecule. Ionic, Atomic, and Molecular Lattices

We have already interpreted the crystal as a limiting case of a very large but unusually regular and symmetrical molecule. This point of view presents only one side of the picture and thus has its limitations. However, it allows us to understand many crystal phenomena without further explanation. It seems clear, for instance, that there can be unlocalized electrons in crystals, just as in the benzene ring molecules (page 415), which are only part of the crystal as a whole. Conversely, there are other electrons, just as the emitting electron of a chromophoric group of a complicated polyatomic molecule (page 402), which we can attribute to a definite group of atoms or even to a definite atom or ion of a crystal. In discussing the optical properties as well as the conductivity of crystals, this difference will prove to be very important.

A parallel can also be drawn between molecules and crystals with respect to the nature of their bonds. Thus we find that in crystals there is a classification according to the type of binding which corresponds to what in molecules we called polar, homopolar, and van der Waals binding. The *ionic crystal lattice* with simple ionic binding corresponds to the polar molecule, and it is apparent that the constituents of the ionic molecule  $\text{Na}^+\text{Cl}^-$  can form only an ionic lattice in the crystal. The alkali-halide crystals are typical representatives of this type of lattice. The *atomic crystal* corresponds to the molecule bound by homopolar forces. Here we must distinguish between two types, those with localized bonds (by electron pairs), and those with nonlocalized binding. The first group is called *valence crystals* and its best known representative is the diamond. The metals and alloys comprise the group with nonlocalized binding.

The *molecular crystal* corresponds to the van der Waals molecule (page 384). It should be called *van der Waals crystal*, because not only molecules, but also atoms can act as crystal constituents (as in the case of the solidified noble gases). If the lattice constituents are molecules

(e.g.,  $H_2$ ,  $Cl_2$ ,  $CO_2$ ,  $CH_4$ ), whose atoms can be bound by polar or homopolar forces, the molecules themselves are held in the lattice by the weak van der Waals forces (page 384). Accordingly, these molecular crystals (just as the noble-gas crystals) loosen up, i.e., melt at very low temperatures; in other words, they require for their decomposition only a small thermal energy. All organic compounds, for example, crystallize in molecular lattices. As a transitional type we should mention the "layer crystal" type in which (to choose at random)  $LiOH$ ,  $CdI_2$ , and  $CrCl_3$  crystallize. In these crystals we have a genuine atomic or ionic binding (homopolar or polar) within the crystal layer, while the individual crystal layers are held together only by van der Waals forces.

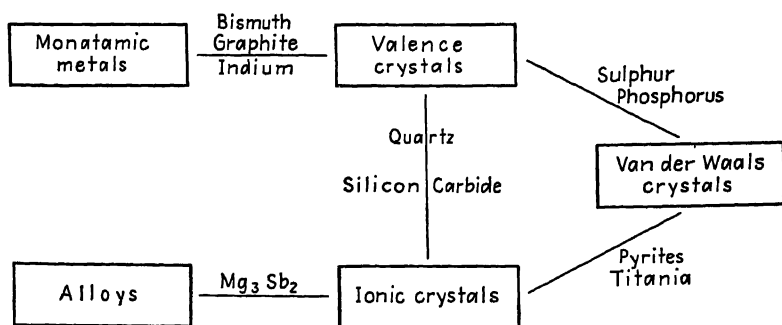


Fig. 198. Schematic diagram of the relation between the different types of solids. (After Seitz.)

A diagram according to Seitz (Fig. 198) presents the relation between the different types of solids in a pictorial manner. The monatomic metals have a definite relationship to the valence crystals, between which the crystals of bismuth, indium, and graphite form transitional cases. The alloys, and according to Pauling also the pure monatomic metals, have some relationship to the polar-bound ionic crystals. Evidence for this can be derived from the transitional case of certain intermetallic compounds, such as  $Mg_3Sb_2$ , with fairly well localized bonds of the ionic type. Furthermore, there are transitions between the valence and ionic crystals, just as in the analogous case of molecular binding (page 414). Thus one can consider the crystals quartz,  $SiO_2$ , and Carborundum,  $SiC$ , as valence crystals with some contribution of ionic binding. Nor are the van der Waals crystals completely isolated, since solid sulphur and phosphorus represent transitional cases between them and the valence crystals and, on the other hand, ice, pyrite, and  $TiO_2$  form transitional cases to the ionic crystals.

Before we consider further the question of lattice binding, at least



for the more important cases of ionic crystals and metals, we must briefly review the geometric arrangement and the determination of the lattice structure of a crystal.

#### 7-4. Crystal Systems and Structure Analysis

We have already explained that because of the symmetric forces acting between the atoms or ions of a solid body (crystal) a stable equilibrium is possible only if the lattice constituents are arranged in a regular geometric pattern. By group-theoretical methods it can be proved that there is a finite but large number (230) of crystal systems which are possible arrangements of the lattice constituents in a crystal. In which one of these systems a particular material crystallizes depends upon the size of the lattice constituents and upon the magnitude and orientation of the forces acting between them. The same material can crystallize in different systems depending upon the type of binding forces; for example carbon can crystallize in the graphite or diamond lattice. The determination of the structure of a crystal, the so-called structure analysis, often is not simple; it is the task of the crystallographer. It is achieved almost exclusively by means of X-ray diffraction following the procedure of von Laue or Debye-Scherrer. We need not go into the details of these methods here. Some remarks on the basis of the method may suffice.

It follows from the periodicity of the lattice structure that for each crystal there is a smallest element, the *elementary cell*, from which the whole crystal can be built by its translatory repetition in the three coordinate directions. In the cubic NaCl crystal the elementary cell (Fig. 199) is a small cube; in most cases the form of the elementary cell is much more complicated. If a coordinate system is constructed whose axes are parallel to the edges of the crystal, then each crystal plane (and thus the whole crystal as determined by the crystal planes) can be characterized by specifying the lengths from the origin to the intersections of every crystal plane with the axes. These lengths are measured in units of the corresponding edges of the elementary cell. The system generally applied for characterizing crystal planes uses not the values from the origin to the intersections of the planes with the axes themselves, but their reciprocals which, when reduced to smallest integral numbers, are called the *Miller indices*. Thus, for example, the

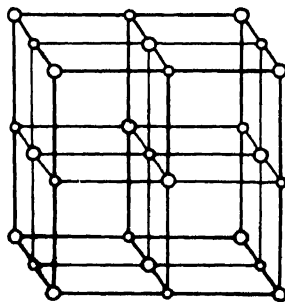


FIG. 199. Unit cell of the rock salt crystal NaCl. The  $\text{Na}^+$  and  $\text{Cl}^-$  ions are indicated by their different size.

(123)-face corresponds to the intersected lengths  $1, \frac{1}{2}, \frac{1}{3}$  or to the smallest integral numbers 6, 3, 2. The (100)-face, on the other hand, is a crystal plane parallel to the  $yz$  plane.

To determine the system of a given crystal, one uses the fact that when X-rays traverse a crystal, they are reflected by the atoms occupying a lattice plane in the crystal, or, in other words, they are diffracted by the lattice points. The relation between the deflection angle  $\alpha$  (to be taken from the Laue pattern), the wavelength  $\lambda$  of the X-rays, and the distance  $d$  between successive lattice planes is given by the Bragg formula

$$2d \sin \alpha = n\lambda \quad n = 1, 2, 3, \dots \quad (7-1)$$

Moreover, we can see immediately from the Laue pattern, Fig. 200, that the symmetry of this particular crystal with respect to the direction

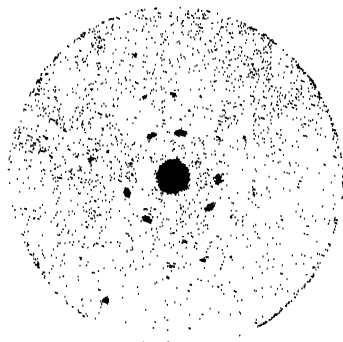


FIG. 200. Laue diagram of a crystal with three-fold symmetry.

of the incident X-rays is a threefold symmetry, i.e., a rotational periodicity of  $120^\circ$ . By taking Laue diagrams, with directions of incidence parallel to the different crystal axes and to the diagonals, a clear picture of the symmetry of the whole crystal can be obtained, whereas the absolute distances between the different lattice planes follow from Debye-Scherrer photographs (page 74) taken with monochromatic X-rays of known wavelength. We have indicated this crystal analysis only quite briefly because it is of no further interest to us

in the atomistic treatment of the properties of solid bodies.

However, the so-called *Fourier analysis*, a theoretically difficult and rather cumbersome quantitative evaluation of X-ray diffraction photographs, taking into account also the intensities of the radiation scattered in different directions, provides not only the arrangement of the crystal constituents in the crystal but also a means of determining the electron distribution effecting the binding between the atoms or ions. Figure 201 shows as an example the electron density in an NaCl crystal in which the lines indicate the places of equal electron density and the numerals, the number of electrons per cube of 1 Å on a side. This method is especially important in investigating the character of the crystal binding in certain cases of transition between one type of binding and another. We have already discussed this Fourier analysis on page 348 in our discussions of the methods of research in molecular physics.

### 7-5. Lattice Energy, Elasticity, Compressibility, and Thermal Expansion of Ionic Crystals

After this survey of the methods of determining crystal structures we shall consider crystal binding, and the phenomena depending upon

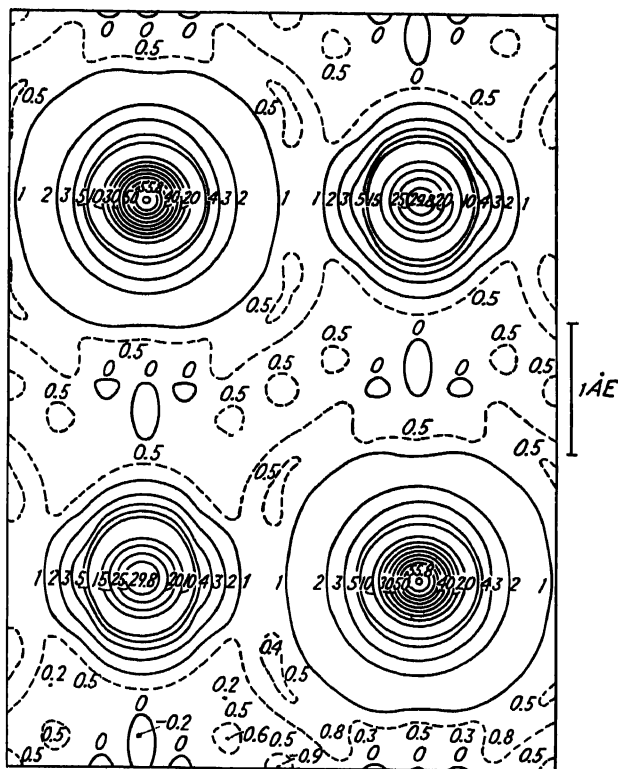


FIG. 201. Distribution of the electron density in a NaCl crystal, determined by Fourier analysis of X-ray data. The numerals indicate the electron density per cube of 1 Å on an edge. (After Grimm, Brill, and coworkers.)

it, more carefully for ionic crystals. If we choose the model of the ideal rock salt crystal then the binding depends on the electrostatic attraction between the singly charged ions  $\text{Na}^+$  and  $\text{Cl}^-$ . Consequently, the potential of this force of attraction between two isolated ions is

$$u(r) = -\frac{e^2}{r} \quad (7-2)$$

However, the two ions are not isolated in space but are in an environment of similar positive sodium and negative chlorine ions. One can

easily see that the nearest neighbors of the ion pair under consideration, because of their opposing charges, tend to reduce the binding. Quite generally, all ions of the environment will contribute to the binding energy positive or negative amounts, which decrease with increasing distance from the ion pair under consideration. The exact computations as carried out by Madelung show that the above expression for the attractive potential must be multiplied by a factor 0.29 in order to take into account the contributions of the other lattice constituents.

Since the two ions cannot penetrate each other, there must be a repulsive force in addition to the attractive force. This repulsive force can be described in a pictorial manner. Any attempt to decrease the distance between two ions, if they touch each other, requires a deformation of their shells, and this deformation requires a considerable amount of energy. For this reason, the potential of the repulsive forces increases with a high power of  $1/r$ . The experimental results can best be described by the ninth power of  $1/r$ . Thus we have for the potential of the forces acting between two ions

$$u(r) = -0.29 \frac{e^2}{r} + \frac{c}{r^9} \quad (7-3)$$

where  $c$  is a constant. In order to determine its value, we make use of the fact that for the equilibrium internuclear distance  $r_0$  the potential energy must be a minimum,

$$\left( \frac{\partial u}{\partial r} \right)_{r=r_0} = 0 \quad (7-4)$$

By performing the differentiation we get

$$c = \frac{0.29}{9} e^2 r_0^8 \quad (7-5)$$

and thus

$$u(r) = -0.29e^2 \left( \frac{1}{r} - \frac{1}{9} \frac{r_0^8}{r^9} \right) \quad (7-6)$$

By plotting the potential  $u(r)$ , just as we have done in molecular physics (page 367), we get the diagram Fig. 202. So far we have only considered the binding between a  $\text{Na}^+$  and one of its neighboring  $\text{Cl}^-$  ions. The term *lattice potential* or *lattice energy* means the total energy per mole of the crystal, i.e., the energy required per mole to form the crystal from widely separated  $\text{Na}^+$  and  $\text{Cl}^-$  ions. To compute the lattice energy, we have to consider that each  $\text{Na}^+$  ion in the lattice has six neighboring  $\text{Cl}^-$  ions, each of which is bound to it by the same  $u(r)$  computed above. Furthermore, 1 mole of a NaCl crystal has  $N_0$   $\text{Na}^+$

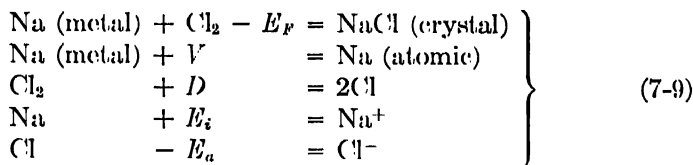
ions from each of which binding forces extend to each of its six neighbors. Thus the lattice energy (referred to a normal crystal with an equilibrium distance  $r_0$  between two consecutive constituents) is

$$U = 6N_0u = -1.74N_0e^2 \frac{8}{9} \frac{1}{r_0} \quad (7-7)$$

This lattice energy is negative because it is released if the crystal is formed from ions which are assumed to exist. If we let the numerical values which refer to the specific case of an NaCl crystal and the constants be contained in a specific constant  $C$ , and if we take into consideration that  $r_0$  is equal to one half of the lattice constant  $d$ , then the energy of all ionic crystals can be expressed by

$$U = C \frac{e^2}{d} \quad (7-8)$$

The lattice energies computed according to this theoretical formula cannot be compared directly with experimental data, because the experimentally measured heat of formation  $E_F$  of the NaCl crystal refers not to the ions which make up the crystal but to metallic sodium and  $\text{Cl}_2$  molecules. However, Born has shown that the lattice energy  $U$  can be computed from the measured heat of formation  $E_F$ , provided some atomic data are known or can be measured. We denote the measured heat of formation of the crystal by  $E_F$ , the heat of vaporization of sodium metal (always referred to 1 mole) by  $V$ , the dissociation energy of the chlorine molecule by  $D$ , the ionization energy of the sodium atom by  $E_i$ , and by  $E_a$  the electron affinity energy, which is set free when an electron is attached to a Cl atom. We can then write the Born cycle in the following way:



From this the lattice energy required to form a NaCl crystal from the  $\text{Na}^+$  and  $\text{Cl}^-$  ions is

$$U = E_F - V - D - E_i + E_a \quad (7-10)$$

Table 15 shows the good agreement of a number of lattice energies

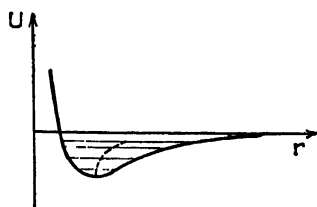


FIG. 202. Interaction potential  $u(r)$  between two crystal constituents. The average internuclear distance is indicated as a function of the vibration energy.

computed from the theoretical formula (7-7) with those determined from experimental data by means of the Born cycle (7-10).

*Table 15. Lattice Energies of Several Alkali-halide Crystals in Cal/Mole. Comparison of the Measured and Computed Values*

Alkali-halide crystal	Experimental value	Computed value
NaCl	183	182
NaBr	170	171
NaI	159	158
KCl	165	162
KBr	154	155
KI	144	144
RbCl	161	155
RbBr	151	155
RbI	141	138

From the curve of the mutual potential  $u(r)$  of two lattice constituents of an ionic crystal (7-6), derived on page 428, a number of important macroscopic properties of crystals can be easily deduced or explained in terms of the atomic theory.

The compressibility is defined as the decrease in volume per unit of pressure. That energy must be applied to reduce the volume, i.e., decrease the equilibrium distance  $r_0$ , can be seen from the potential curve. Moreover, it can be seen immediately that the compressibility of a crystal is the smaller, the steeper the potential curve rises from the minimum toward smaller internuclear distances. The force  $F$  necessary to reduce the distance can be computed by differentiating  $u(r)$  with respect to  $r$ . Thus the work per mole required to reduce all distances from  $r_0$  to  $r$  is

$$A = 6N_0 \int_{r_0}^r F dr = \int_{r_0}^r p dv = \frac{P}{2} \Delta v \quad (7-11)$$

Here  $p$  is the instantaneous pressure which increases with the compression and  $P$  is the final pressure corresponding to the total change in volume,  $v$ . Here the function  $p(v)$  has been approximated by a linear function. Carrying through the computation results in an expression for the compressibility,

$$K = c' \frac{d^4}{e^2} \quad (7-12)$$

The compressibility thus increases with the fourth power of the lattice constant  $d$ , a plausible result since the compressibility must be the smaller the more densely the lattice components are packed in the normal lattice.

If the compressing force is removed, the equilibrium distance  $r_0$  is restored, since without external forces equilibrium can only exist if all distances correspond to the minimum of the potential curve. *Consequently, the crystal is elastic.* From this fact it follows that the crystal can vibrate, i.e., periodically increase and decrease the distance between the lattice points. We shall consider this important property in some detail in the section following the next.

The last important conclusion which can be drawn from the potential curve concerns the thermal expansion of crystals and, in general, of all solid bodies. The vibration energy of the crystal and thus the amplitude of the crystal vibrations increases with increasing temperature. Because of the asymmetry of the potential curve (anharmonicity of the lattice vibration) the average distance between two lattice points, shown in Fig. 202, increases with increasing amplitude of the vibration. Consequently, the volume of the crystal increases with increasing temperature. *Thus the fundamental thermal property of all known solid bodies, their expansion with increasing temperature, has a very simple explanation from the atomic theory. In the last analysis it depends upon the asymmetry of the potential curve,  $u(r)$ .*

We have limited ourselves in this section to a discussion of ionic crystals because their properties can be accounted for quantitatively. Qualitatively, the behavior of valence crystals, such as diamond, with respect to the potential curve and the conclusions drawn from it, such as compressibility, thermal expansion, etc., is very similar to that of the ionic crystals. Phenomenologically, valence crystals are characterized by their great hardness and good electrical insulating properties. The latter is due to the fact that electrolytic conductivity, which will be considered in detail for ionic crystals on page 457, is not possible because the constituents are atoms. Electronic conductivity, on the other hand, is impossible because, due to the strongly localized valence binding, electrons are so strongly bound that there are no conduction electrons available.

## 7-6. Review of the Binding Forces and Properties of Metals

Now that we have considered the properties of insulating crystals with strongly localized bonds, we shall briefly discuss the binding and the atomic properties of the other extreme of solids, the metals. They are characterized by a lack of localized binding. Such a review of the properties of metals appears to be the more appropriate as metals have to be regarded, from a practical technical standpoint, as one of the most important groups of solids. Moreover, many of the properties of metals, their modifications (different forms of lattices in different

temperature regions), and their alloys were first explained by atomic physics which also threw new light on the problem of the technological treatment of metals (e.g., hardening and tempering).

In the course of this chapter we shall consider and explain the characteristic differences between insulators and metals from different points of view. At this time it is sufficient to conclude from their high electrical and thermal conductivity, that the outermost electrons of the atoms forming the metal (in general, one per atom) are not tightly bound to their atoms as they are in most insulators. On the contrary, they must have a high mobility within the metal just as do the molecules of a gas. This is why we often speak of the electron gas in the metal, a degenerate electron gas, according to page 239. These external atomic electrons which, as valence electrons, are also responsible for binding (page 409), are thus by no means localized. According to page 415, the last six valence electrons of the  $C_6H_6$  molecule are distributed, as nonlocalized electrons, among all six C atoms of the benzene ring, and thus participate in its binding, though they do not form electron pair bonds. In a similar way the nonlocalized conduction electrons of a metal must be responsible for the binding in the metallic state (metal crystal). However, again there is no mutual saturation of electron pairs to form localized bonds. In contrast to homopolar binding, *metallic binding thus does not show the phenomenon of saturation*. Whereas two homopolarly bound H atoms cannot bind an additional H atom, in a metal each additional atom is bound with the same force independently of the number of atoms already bound in the crystal.

We know from experience that the elements on the left side of the periodic table exhibit metallic character, whereas localized binding predominates in crystals formed from the elements on the right side of the table, and it is here that we find the insulating crystals. According to Hume-Rothery, the reason for this distinction is that metals generally crystallize in a body-centered or face-centered cubic lattice or in a hexagonal, close-packed structure. On the other hand, the elements on the right side of the table belong to very different lattice types in which the coordination number (the number of surrounding atoms at the same smallest distance from one atom) is always equal to the valence of the atom under consideration. According to Fröhlich this difference depends upon the repulsive Coulomb forces between each one and all the outer electrons. Sodium and chlorine, for example, each have the valence 1. Sodium has only one outer electron, whereas chlorine, which lies on the right side of the periodic table, has seven. The binding effect of the one electron pair in the Na—Na and Cl—Cl pairs therefore is different, because in Cl—Cl we have an additional repulsion between the two times



six excess electrons. Such a repulsion is not present in the case of Na—Na. Nevertheless, in the case of the diatomic  $\text{Cl}_2$  molecule, there is a strong homopolar binding because the exchange forces (page 412) produce a unilateral charge displacement (and thus an attraction between the atoms) which compensates for the repulsion due to the other electrons. However, such an asymmetry of the electron configuration is incompatible with the high symmetry (coordination number) of the crystal lattice. Thus, because there is no Coulomb repulsion, sodium atoms can be bound by metallic binding in the Na-metal lattice, whereas for Cl atoms such a symmetrical binding is not possible. The Cl atoms can form  $\text{Cl}_2$  molecules by unilateral deformation and reduction of their internuclear distance. These molecules can then be bound in a crystal lattice by the relatively small van der Waals forces (page 386). This contrast, which was explained here as an example for Na and Cl, *holds in general for all elements on the left or right side of the periodic table: on the left side we have metallic binding with high symmetry and coordination number, on the right complicated lattices strongly influenced by the valence and always without electrical conductivity.*

We may also mention that other important properties of all solids, such as compressibility and plasticity, are related to these differences. Both decrease in going from the left to the right, as a result of the increasing effect of the repulsive action of the outer electrons.

Metallic binding is thus characterized by the lack of valence saturation, and therefore by a high symmetry (equally strong binding of all equally distant neighbors independent of the valence of the metallic constituents). We shall try to understand the reason for this behavior of metals. Because of the high metallic density, the distance from atom to atom in the lattice is smaller with reference to the atomic radii than in most insulators. We have, therefore, a strong overlapping of the electron shells of the original atoms which form the metal lattice. This overlapping means, according to the exchange theory (page 233), that there is a high exchange probability of the outermost electrons between neighboring atoms, and thus a large exchange binding energy. However, Dehlinger and Pauling have proved that we may also look at the problem from another point of view. We may speak of a normal formation of electron pairs, and thus emphasize the relation between metallic binding and normal valence binding. However, we have then to consider these electron pair bonds between one atom and its neighbors as rotating, so that over an average of time we have again a symmetrical (nonlocalized) distribution of the electrons which cause the binding. It can be shown thermodynamically that this state with rotating bonds can be more stable at higher temperatures than the states of unilateral

homopolar or polar binding. This case seems to be realized in many metal alloy lattices which have higher coordination numbers at higher temperatures than they do at lower temperatures.

Pauling has extended this picture by the additional assumption (whose correctness he was able to show thermodynamically), that ionic states also contribute to the normal metallic bonds. Thus a monovalent alkali atom in a metal occasionally should have two valence electrons and thus, as a negative ion, should be able to bind simultaneously two of its neighbors.

In general it can be said that in metals the nonselective metallic binding predominates and determines the metallic character. Nevertheless, the individual properties of the atoms, especially their spin, orbital momentum, and the valence determined by them, play a certain role. Thus, in addition to the general metallic binding, there is a certain amount of heteropolar or homopolar binding. Consequently, *a metal always crystallizes in the lattice of highest symmetry which is compatible with the special properties of the atomic lattice components.* For example, in the case of the alkalis, it follows from the diamagnetism of the metal that the spin momenta of the outer electrons of the atoms are alternately antiparallel. For this case, quantum-mechanical computations (by means of the exchange integral, pages 231 and 232) show that the body-centered lattice, in which each atom is surrounded by eight atoms of oppositely directed spin, is more stable than the lattice expected for pure metallic binding, in which each atom is surrounded by 12 others at equal distances, half of which have one spin direction and half the other. The energetically necessary antiparallel spin directions of the alkali electrons, which are empirically demonstrated in the diamagnetism of the alkali metals, thus are responsible for the fact that, not the lattice corresponding to pure metallic binding with a coordination number 12, but the body-centered lattice with coordination number 8, has the lower state of potential energy. Any other structure and any other spin distribution has a lower stability. The body-centered lattice which is expected theoretically has actually been found for the alkali metals by X-ray experiments.

In a similar manner the empirically established lattice structure of metals and alloys, which in turn accounts for many other properties of metals, can be explained in great detail from the point of view of atomic physics. An especially good example is presented by the different modifications of iron, i.e., the different stable lattice structures at different temperatures, which are only partly ferromagnetic. These can be understood by a detailed discussion of the effect of exchange energy (page 451). The second technically significant example of the co-

operation of metallic and homopolar or polar binding is that of the alloys, i.e., mixed crystals of different metals. In general, substitutional mixed crystals are observed in alloys (if the difference in the atomic radii is not too great), in which both metals occur in arbitrary proportions. However, it follows from sudden changes of the electrical resistance as well as of the magnetic properties, that for certain rational proportions of the components (corresponding, for example, to the compounds AuCu, AuCu<sub>3</sub>, FeCo, or Ni<sub>3</sub>Fe) especially stable so-called superstructures occur which provide further evidence that the general metallic binding is superimposed by some localized binding which exhibits saturation characteristics. These localized valence forces which cause an ordered structure of alloys (e.g., Cu ions in the center and Zn ions at the corners of the elementary cubes in  $\beta$ -brass) are very weak forces. Evidence for this is the observation that at temperatures of only a few hundred degrees this ordered structure disappears and is replaced by a disordered lattice with a probability distribution of Cu and Zn ions.

For the sake of completeness, we mention that, besides the substitutional alloys, there exist also interstitial alloys, in which small ions like H, C, or N occupy interstitial places (see page 454). The systems palladium-hydrogen and iron-carbon are examples of such interstitial alloys.

Among the important properties of metals, to be explained by atomic physics, we finally mention ferromagnetism and superconductivity. The former will be discussed after we have learned more about the behavior of electrons in metal crystals, whereas superconductivity cannot as yet be satisfactorily explained in spite of some promising approaches.

### 7-7. Crystal Vibrations and the Determination of Their Frequencies from Infrared and Raman Spectra

Because there is a quasi-elastic binding force effective between the lattice constituents of a crystal, vibrations between them are possible just as they are between the nuclei of molecules (page 365). The asymmetrical potential curve  $u(r)$  (Fig. 202) is evidence that we are dealing with anharmonic vibrations. We have already used this fact in explaining the thermal expansion of solids.

We have to distinguish between the so-called *internal vibrations* and the *lattice vibrations proper* of a crystal. Internal vibrations occur particularly in molecular crystals (page 423), in which real molecular vibrations are possible within the molecules which form the lattice constituents. These molecular vibrations may be disturbed more or less by the crystal lattice, but have nothing else to do with the whole crystal.

In the lattice vibrations proper, on the other hand, the whole crystal participates. To illustrate these two types of vibrations we shall use the calcite crystal  $\text{CaCO}_3$ . In this crystal the lattice components are the ions  $\text{Ca}^{++}$  and  $(\text{CO}_3)^{-}$ . The vibrations of all ions of one kind with respect to the ions of the other kind are called the external or lattice vibrations, because the whole lattice participates in this vibration. On the other hand, there are internal vibrations within the  $\text{CO}_3$  group (e.g., the O atoms vibrate with respect to the C atom) which have nothing to do with the calcite crystal as a whole. Therefore they occur with the same frequencies in other crystals containing the  $\text{CO}_3$  group, such as  $\text{FeCO}_3$  or  $\text{MgCO}_3$ .

We have to distinguish two types of lattice vibrations. In the first type, all identical lattice constituents vibrate in phase (*optical vibration*), whereas in the second type the vibrations are not in phase (so-called *acoustic vibrations*). In the first case we have, at least in ionic crystals, a vibration of all positive ions with respect to all negative ions, i.e., a vibration of the whole  $\text{Na}^+$  lattice with respect to the  $\text{Cl}^-$  lattice. As a result of the corresponding change of the dipole moment this vibration must appear in the infrared absorption spectrum. For this reason we speak of optical vibrations. In the other extreme case the vibration progresses in the crystal as a longitudinal wave. Since in this case similar atoms do not vibrate in phase, the changes in the dipole moments largely compensate each other. Therefore these vibrations do not appear in the infrared absorption spectrum and are called acoustic vibrations since longitudinal waves of this type are known from acoustics. They are less important for crystal physics than the optical or fundamental vibrations of the lattice.

The frequency of the fundamental vibration associated with a change of the dipole moment (which may originate from the polarization of the vibrating atoms) can be taken from the infrared absorption spectrum of the crystal. The absorption coefficient of the corresponding wavelengths is very large, because all lattice ions participate in this vibration. Absorption measurements therefore are made with very thin crystal laminae. A more elegant and simpler method, although not so exact, is the reststrahlen method of Rubens. Because of the relation between absorption and reflection, the crystal reflects strongly those wavelengths which correspond to the fundamental vibration, whereas the other wavelengths are mainly transmitted by the crystal. If a light ray is reflected back and forth many times between two plates of the crystal under investigation, then all wavelengths except those of the fundamental vibration are gradually eliminated by transmission. The residual rays, or reststrahlen, remaining after many reflections thus

consist essentially of the fundamental vibrations whose wavelengths can easily be measured. Occasionally it happens that, as a result of special symmetry relations which exist, for example, in fluorspar  $\text{CaF}_2$ , the fundamental vibration is not associated with a change of the total dipole moment and thus it is optically inactive. In that case it cannot be found in the absorption spectrum.

If the polarizability of the lattice constituents is changed by such optically inactive vibrations, as is almost always the case, the frequencies of the vibration under consideration can be determined by Raman effect measurements as in the corresponding case in molecular physics (page 352). This implies the advantage of working in the visible or ultraviolet region, so that optical spectrographs with their higher resolving power and dispersion can be used. Ultraviolet and Raman investigations again supplement each other so that there is no basic difficulty in determining the frequencies of the fundamental lattice vibrations.

### 7-8. Electron Configuration and Electron Transition Spectra in Crystals. The Energy Band Model

After the discussion of the types of solids, crystal structure, and crystal vibrations, we turn to the phenomena depending upon the electrons in crystals. We investigate first the electron configuration. The problem of the electron configuration in crystals and of the spectra resulting from electron transitions can be approached from very different points of view, all leading to the same result.

We consider first the transition from an undisturbed atom to one which is strongly disturbed by its environment (atom in a plasma, page 187), and finally to one which is incorporated in a crystal as a greatly disturbed lattice atom. By this very pictorial approach, we arrive at an essentially correct picture. Undisturbed atoms have, according to page 106, the sharp energy levels of the bound and the continuous energy regions of the free (ionized) electrons. In the spectra of plasma atoms which are disturbed by the microfields of their environment, the broadening of the higher energy states of the electrons as a result of the perturbation is already very noticeable (page 187). For the highest energy states just below the ionization limit, it can no longer be decided whether they still are to be considered as discrete energy states of the bound electrons or as continuous energy regions of the free electrons (reduction of the effective ionization potential). Finally, if we go to the lattice constituents in a crystal, we expect a different behavior, aside from the increasing magnitude of the perturbation and a broadening of the energy states, only insofar as *there is no longer a completely*

free electron. An electron released from its original atom will now move "quasi-freely" throughout the crystal in the periodic potential field of all lattice ions. We have thus arrived, in a pictorial way, at the correct picture of the electrons in a crystal. It is evident and clear that by incorporating the atom in a crystal lattice, we do not influence the inner electrons which are responsible for X-rays. These innermost electrons remain strongly bound to their respective nuclei; their energy states are practically undisturbed and therefore remain sharp. Evidence for this is the sharpness of the X-ray lines emitted by metallic anticathodes (see Fig. 67, page 119). The perturbation and thus the breadth of the energy levels of electrons increases greatly with increasing principal quantum number  $n$ . In the case of the optical energy levels, which in general are not occupied by electrons, this width reaches several electron volts, so that one can readily speak of energy bands of the electrons which are now already quasi-free. We shall see that these quasi-free electrons are responsible for the electronic conductivity of the metals (page 443). For the highest energy levels the band width actually becomes so great that there is an overlapping of the different bands. We shall discuss this when we take up the theory of metallic conductivity. An energy band diagram is shown in Fig. 203.

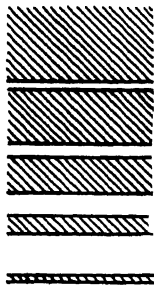


FIG. 203. Energy-band diagram of the electrons in a crystal (schematic).

It is interesting to show how one can arrive at the same result by an entirely different approach. For that reason we mention two other concepts. They are less pictorial but they have the advantage that they permit an exact computation of the behavior of the crystal electrons.

One approach is based on the concept of resonance or exchange splitting which was discussed in Chap. 4. Quantum mechanics shows (page 230), that in the case of two coupled atomic systems of equal energy their common energy state splits into two energy states, as a result of this energy resonance. The separation of these new states is the larger, the stronger the coupling between the two systems. This case exists in crystals where in principle an exchange of electrons by any two lattice constituents is possible because of their complete identity. Thus if the crystal consists of  $N$  atoms, then as a result of the possibility of exchange of each electron with the  $N-1$  analogous electrons, each energy state of the atoms in the crystal splits up into  $N$  energy levels. We distinguish these  $N$  energy levels of every energy band by quantum numbers  $k$ . This  $k$  is closely related to the momentum of the particular electron and thus corresponds to the orbital quantum number  $l$ , intro-

duced on page 130 for the bound atomic electrons. According to the Pauli principle, each  $k$  state can be occupied by two electrons of opposite spin. *The magnitude of the splitting, and thus the width of the band resulting from the  $N$  levels, depends on the degree of coupling, i.e. on the exchange probability of the electrons [Eq. (4-121, page 231)].* We discuss this electron exchange by means of the potential curve diagram Fig. 204. In Fig. 105 (page 187) we had represented the course of the potential near a hydrogen nucleus together with the energy states of the electron which moves in this potential field. *In a solid, because of the geometrically regular arrangement of the atoms, we have a three-dimensional periodic*

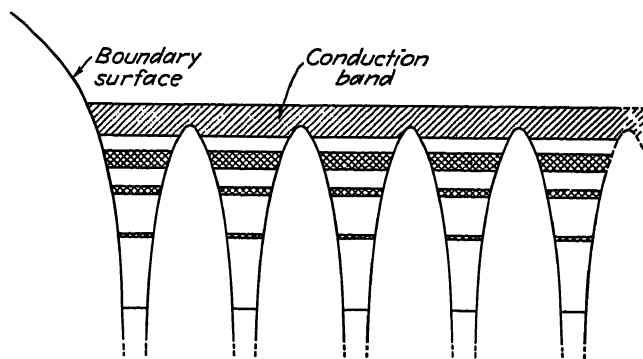


FIG. 204. Potential field in a crystal with electron states, excitation bands, and ionization (conduction) band (schematic).

*repetition of potential wells and potential walls, which are indicated in Fig. 204 with the corresponding energy bands. The innermost electrons of the crystal atoms lie almost completely within the potential troughs of their ions. An exchange of these electrons with corresponding electrons of other atoms is impossible classically. Because of the tunnel effect (page 234), it is not quite impossible according to quantum mechanics, but it is very improbable because of the height of the potential wall which must be penetrated. Because of the small exchange probability, the splitting of the energy levels, and thus the breadth of the energy bands of the innermost electrons, in agreement with our first picture and with experience, is very small. Conversely, the exchange probability and the breadth of the energy bands resulting from the energy levels is very great for the quasi-free external electrons. Thus our picture is in agreement in its results with our first grossly pictorial derivation. However, it shows that the energy bands consist theoretically of discrete energy levels and are practically continuous only because of the large number of these levels. This number per mole of the crystal equals the number of lattice atoms and thus amounts to  $6 \times 10^{23}$ .*

In our third method of representation we start from the quasi-free electrons in the crystal and we shall show that there must be certain forbidden energy regions, the regions lying between the energy bands. To do this we consider electrons moving in the lattice according to wave mechanics as progressive electron waves whose wavelength  $\lambda$  (page 197) is

$$\lambda = \frac{h}{mv} = \frac{h}{\sqrt{2mE}} \quad (7-13)$$

if  $E$  is the energy of the electrons. Now it can be shown geometrically that only for certain definite wavelengths  $\lambda$  or certain energy values  $E$  the electron wave can penetrate the crystal. For different values of  $E$  or  $\lambda$  the electron wave is reflected by the lattice planes in such a way that it is extinguished by interference. The quantitative treatment of these considerations, which are only indicated here, leads naturally to the same results as the other two methods of approach. For this exact computation an expression for the three-dimensional periodic lattice potential is sought in which the electrons move. This expression is inserted in the Schrödinger equation (4-42), page 207. This equation has for solutions the eigenfunctions which describe the behavior of the electrons in the crystal, only for certain energy regions, i.e., for exactly the energy bands of our diagram Fig. 203.

In order to prevent a possible misunderstanding we want to emphasize one point. The broad energy bands of solids are pure electron energy bands and cannot be regarded as smeared bands of band systems of macromolecules. In molecules (page 364), electronic motion and atomic vibration are closely coupled with one another because an excitation of an electron produces in general a change in the molecular binding and as a result produces a vibrational excitation (Franck-Condon principle, page 370). Thus in a molecule an electron transition generally is associated with a change of vibration (rotation). A whole band system results from *one* electron transition because of the large number of vibrational and rotational energy changes coupled with it. If we disregard the internal vibrations in the molecular crystals, *crystal vibrations affect, in general, the whole crystal lattice. It seems evident that lattice vibrations are less influenced by the excitation of individual electrons.* Thus in a crystal the coupling between the motion of electrons and the vibration is very small. To an electron transition from one energy band to another in a crystal there corresponds a spectrum which might be regarded as a broadened spectral line and is usually called a spectral band. It should be kept in mind, however, that it has nothing to do with a molecular band or band system, because in the case of a spectral



band of a crystal the characteristic property of a band spectrum, the participation of the vibration and rotation energy, is missing.

We return now to the energy bands of the crystal electrons. The distribution of the  $k$  levels (see page 438) over each energy band can be taken from the spectrum: its density is greatest in the center of the band and it decreases to zero at the edges of the band. The intensity distribution of a spectrum resulting from a transition from a sharp (inner) level to an energy band thus is similar to that of a broadened spectral line (the bell shaped curve, Fig. 103).

For the selection rule for optical transitions between energy bands it follows from the theory that transitions between all the different energy bands (different  $n$ 's) are allowed, but with the important restriction that in optical transitions the quantum number  $k$  must not change. Since  $k$  is related to the momentum of the electron, this is essentially a consequence of the law of conservation of momentum. Surprisingly, the enumeration of  $k$  in the energy bands is not the same. If, as assumed in Fig. 205a,  $k$  increases from the lower to the upper edge of the energy band, then the short- and longwave optical transitions are shown by the arrows in Fig. 205a. The width of the corresponding spectral band is equal to the difference in the widths of the two combining energy bands. However, if  $k$  increases in one band from the lower to the upper edge of the band, and in the other band from the upper to the lower edge, then the possible short- and longwave transitions are shown by the arrows in Fig. 205b. In this case the width of the spectral band is given by the sum of the widths of the energy bands. In the case of metals, for example, both cases occur. We shall see on page 470 how important the  $k$ -selection rule is for understanding the luminescence of crystals.

A very nice experimental proof for the width and arrangement of the nonoccupied optical energy bands can be taken from X-ray absorption spectra. According to page 123, the longwave limit of the X-ray absorption continua originates from raising an electron to an unoccupied energy level near the ionization limit of the atom. The longwave structure of an X-ray absorption edge thus reflects directly the energy band structure in the crystal. This has been found to be in best agreement with the theoretical results developed above.

There is little to say about the optical spectra of solids associated with electron transitions so long as we speak of pure and ideal crystals. Absorption spectra usually result from transitions of electrons from the

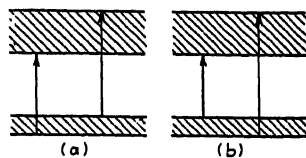


FIG. 205. Different cases of allowed transitions between two energy bands.

highest occupied energy band into one of the higher unoccupied bands. Depending on the different magnitudes of the interaction of the electrons, i.e. their exchange probability, we find broad continuous absorption bands in the ultraviolet and partly in the visible spectral regions for the metals, and very narrow bands, almost lines, for the crystals of the rare earths. This last result agrees well with our pictorial concept: the spectra of the rare earths result from transitions of electrons within the inner  $4f$  shell which is shielded by the  $5$ -quantum electrons (see Fig. 99) against external influences. These energy bands, consequently, are broadened only slightly. The large number of these transitions is due to the fact that the selection rules which are valid for isolated atoms are made invalid by the external influence of the interatomic electric fields in the crystal (see page 129).

Emission spectra of pure, ideal crystals can result from transitions of excited electrons from a normally unoccupied band to holes of a lower band, but these are very improbable, since the requirement of equal  $k$ -quantum numbers for the excited electron and the hole in the nearly filled lower band is almost never fulfilled. We shall become acquainted with the details of this on page 470).

### **7-9. Completely and Incompletely Occupied Energy Bands in Crystals. Electronic Conductors and Insulators According to the Energy Band Model**

We can use the energy band model of the crystal electrons to draw conclusions about the electronic conductivity of crystals, conclusions which are as pictorially clear as they are important. Our aim is to give an atomistic explanation of the fundamental difference between metals and insulators. What are the conditions which determine whether a crystal is a good electronic conductor or an insulator?

In order that electric charge can be transported through a crystal by its electrons it is necessary that, as a result of the movement of electrons in the electric field to the positive pole, there be an excess of electrons on the positive side and a deficit on the negative side. Now if the highest band occupied by electrons is completely filled, and if we disregard the possibility of an electron jumping to a higher band (since this requires very high temperatures or optical excitation), then the accumulation of an excess of electrons on one side of the crystal is not possible. The migration of electrons toward one side of a crystal through a fully occupied band is possible only if simultaneously an equal number of electrons migrates toward the other side. One can understand that a unidirectional migration of electrons in a filled band is impossible, because the electrons would have to take energy from the field if they

are accelerated by it. However, there are no energy states available for these electrons of higher energy in a filled band. Thus, *electron conduction is not possible in crystals with completely filled energy bands. These crystals therefore are called electronic insulators.*

On the other hand, *if in a crystal the highest energy band which contains electrons is not completely filled, then, according to the above statements, unidirectional migration of electrons is possible. Consequently, metallic conductors are characterized by the property that the highest of the occupied energy bands is incompletely filled.*

How does this theoretical result agree with our empirical knowledge of the highest occupied energy bands of the metals? According to our discussion on page 439, each energy band of a crystal consisting of  $N$  atoms has "room" for  $2N$  electrons, i.e., for two electrons per atom. According to the Pauli principle (page 169), each  $k$  state can be occupied by two electrons with oppositely directed spin. We also know that the highest energy band of a crystal is occupied by the valence electrons of the constituent atoms. Consequently, it is clear that the monovalent metals are electrical conductors, because they have only one electron per atom in the highest energy band, which thus is only half filled. However, according to this simplest concept, bivalent metals should be insulators, in contrast to experience. Now in metals the interaction of the electrons which determines the width of their energy bands (and which, according to page 432, also causes the metallic binding of the atoms in the crystal) often is so large that the upper energy bands of the metals partially overlap. Since the state to which the system tends is always that of the smallest potential energy, the  $2N$  valence electrons will distribute themselves in the manner shown in Fig. 206. In this figure the bands are shown one beside the other. Instead of completely filling one band and leaving the other unfilled, the  $2N$  valence electrons thus fill both bands up to the same level. *As a result of this overlapping of the higher bands in a crystal with large interaction between the electrons, the highest occupied bands of the bivalent metals are not completely filled and these solids thus show the electronic conductivity which is characteristic for metals.*



FIG. 206. Overlapping of the highest occupied energy bands of a bivalent metal and occupation of the two bands by electrons. For clarity the two bands are shown side by side.

That the situation presented here is not only a hypothesis but that it corresponds exactly to reality has been shown by a very beautiful experiment of Skinner. The experiment concerns the X-ray emission spectra of various metals. We know from page 438 that the lower electronic states in a crystal can be regarded as sharp levels because of

the small disturbance of the innermost electrons of the atoms by the lattice. The X-ray lines of lighter elements which are emitted after ionizing an electron of the  $K$  shell (page 117), correspond to transitions from the upper broad energy band (so-called conduction band) to the  $K$  level which is very sharp in comparison. The width and intensity distribution of the emitted X-ray band thus gives us a direct picture of the width of the upper energy band and its electron population. Densitometer curves of these X-ray emission bands are shown in Fig. 207 for the monovalent lithium metal and the bivalent magnesium metal. *Only half of a bell-shaped curve is clearly visible in the first case, and this is*

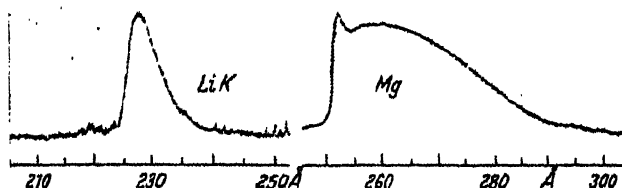


FIG. 207. Densitometer recordings of the X-ray emission bands of a monovalent (Li) and a bivalent metal (Mg), as evidence for the correctness of the band arrangement in Mg according to Fig. 206. (After Skinner.)

*pictorial evidence that the upper band of lithium is only half-filled with electrons. In the case of the bivalent magnesium, on the other hand, it can be easily recognized that the X-ray emission band originates from the transitions from two partially occupied, overlapping energy bands to the same lower state. This is in complete agreement with our discussion given above. In a similar way the energy bands of the alkali-halide crystals and those of a number of oxide crystals have also been studied and attributed to definite electron transitions.*

We conclude with a remark about the temperature dependence of the electronic conductivity. In the ideal nonvibrating crystal, the coupling between the electrons in the conduction band and the lattice built up of their ions is very small and the conductivity is correspondingly large. The lattice vibrations tend to reduce the electron mobility because the electrons collide more often with ions and are scattered irregularly. As the amplitude of the lattice vibrations increases with increasing temperature, the migration of the electrons through the crystal is inhibited to an increasing extent. The electric resistance of the metals increases with increasing temperature, because the electrons must "work their way" through a potential field which varies with respect to space and time. The empirical law according to which the

electronic conductivity at low temperatures is proportional to  $T^{-5}$  thus appears to be explainable in a qualitative way. On the other hand, a conduction phenomenon which cannot be satisfactorily explained from the theoretical point of view is that of supraconductivity. This is the remarkable effect that a large number of metals and metallic compounds seem to lose their electrical resistance completely in the vicinity of the absolute zero of temperature. According to Heisenberg, this phenomenon seems to be associated with a lattice-like arrangement (produced by electrostatic repulsive forces) of those few high-energy conduction electrons which are not yet "frozen" in their lowest possible energy states.

### 7-10. The Potential Well Model of a Metal. Work Function, Thermionic Electron Emission, Field Emission, Contact Potential

In our discussion of solids and particularly of metals from the atomistic viewpoint we have neglected surface effects. We shall now interest ourselves in those phenomena, and especially in the passage of electrons through the surface layer of a crystal into free space (vacuum or air) or into another metal.

In this respect we must *think of the electrons in the conduction band as bound not to the individual lattice ions but to the crystal as a whole and that, consequently, for an electron to escape out of the crystal, energy must be expended against this binding force.* The energy required to release an electron from a metal into free space is called the work function  $\Phi$ .

Thus the potential energy of an electron is greater in free space by the amount  $\Phi$  than it is for the electron in the crystal. We can represent this behavior by a potential curve in the customary manner and thus obtain the potential well model of a metal (Fig. 208). In this representation we are concerned only with the quasi-free conduction electrons. Only their energy levels are, therefore, indicated in Fig. 208. We see that the potential of an electron at rest in the interior of the metal (lower edge of the conduction band) is smaller by an amount  $\Phi_0$  than it is outside the metal. However, this potential difference is not the effective work function for the release of an electron from the metal. The effective work function  $\Phi$  is much smaller than  $\Phi_0$  because the energy levels within the potential well are filled by electrons up to a height  $\delta$ . According to Fermi statistics (page 238) electrons in the highest state already have a zero-point energy equal to  $\delta$ . The effective work function for those electrons which are most easily released from the metal therefore is

$$\Phi = \Phi_0 - \delta \quad (7-14)$$

The behavior described here holds exactly for absolute zero, whereas at higher temperatures there are always a certain number of electrons in higher, formerly unfilled states.

The potential well model gives us a very pictorial explanation for the emission of electrons by an incandescent metal. It is clear that electrons with kinetic energies  $E_k > \Phi$  can leave a metal spontaneously; in the classical picture they evaporate. In 1908, Richardson deduced from measurements a formula for the temperature dependence of the saturation current  $i$  which is obtained when the released electrons are drawn off by a sufficiently high electric field. His result is the famous Richardson equation,

$$i = AT^2 e^{-\Phi/kT} \quad (7-15)$$



FIG. 208. Potential-well model of a metal.  $\Phi$  = effective work function,  $\phi_0$  = work function for electrons in the lowest energy state of the conduction band,  $\delta$  = Fermi surface = zero-point energy of the electrons in the highest occupied electron levels.

$A$  is a constant of the metal which is theoretically equal to 60 if the current density is computed in amperes per square centimeter of the incandescent surface. However, for the large number of materials investigated, the value of  $A$  varies between  $10^{-3}$  and  $10^4$ , if the value of the normal work function is used for  $\Phi$ . This discrepancy is explained by the fact that actually the work function  $\Phi$  of different materials

changes in different ways with the temperature. However, customarily the measurements are represented by the expression (7-15) with constant  $\Phi$ , but different  $A$ . The Richardson equation has been derived by Schottky, Dushman, and others from the classical theory, before the advent of the quantum theory, on the basis of a normal evaporation of the electrons from the heated metal.

Compared to the exponential term  $e^{-\Phi/kT}$ , where  $\Phi$  is the effective work function, the term  $T^2$  is rather unimportant. It seems clear that the electron current has to be proportional to the term  $e^{-\Phi/kT}$ , because this quantity gives the fraction of the electrons which have at the temperature  $T$  the energy  $\Phi$ . The quantum-mechanical derivation from the electron theory of metals leads to exactly the same equation.

Since the work function  $\Phi$  occurs in the exponent of Richardson's equation, a slight reduction in the work function causes a large increase in the number of emitted electrons. The effective work functions for several technologically important metals are listed in Table 16.

Table 16. The Work Function of Several Metals

Metal.....	Cs	Ba	Th	W	Pt
$\Phi$ electron (volts)...	1.8	2.1	3.35	4.5	6.2

We see that the variation in the work function corresponds to that of the ionization potential of the atoms forming the metal. Cesium, for example, has the smallest work function and likewise the smallest ionization potential of all elements (see Table 2, page 30). This relation might be expected from the quantum theory, since *the work function* (actually not  $\Phi$ , but  $\Phi_0$ ) *can be regarded as the ionization energy of the atom when influenced by the crystal lattice.*

The strong dependence of the thermionic emission on the work function is used widely in technical applications. Since the cathode cannot be made of pure cesium, barium, or thorium for various technological reasons, tungsten, which can withstand high temperatures, is used. However, its work function is lowered by depositing on it a layer of cesium or thorium, monatomic, if possible, which results in a large increase of the thermionic electron emission. In the first method, the tungsten cathode is operated in a rarefied cesium vapor atmosphere, by which a monatomic cesium layer is formed on the tungsten. To produce the thorium coating, in the manufacturing process the tungsten is mixed with thorium oxide and this is reduced by heating so that the thorium diffuses toward the surface, thus producing a monatomic thorium layer on the tungsten. This *thoriated tungsten cathode* has the great advantage that it can be used in high vacuum. Surprisingly, by properly coating the tungsten with Cs or Th atoms, a work function is obtained which is only slightly higher than 0.7 or 2.6 ev, respectively, and thus is noticeably *less* than that of pure metallic cesium or thorium. This can be explained by the atomic theory and is confirmed by experiment as being due to the fact that the Cs and Th atoms adsorbed on the tungsten surface are polarized with the positive pole of the resulting dipole directed toward the tungsten surface. In this way the positive charge is stronger and the negative electron cloud less tightly bound than for symmetrical binding. The separation of an electron from this type of adsorbed atoms therefore requires an amount of energy smaller than the normal effective work function. This phenomenon is an excellent example of the atomistic explanation of even complicated details of a technically important effect.

Another application of the model of the potential well is the field emission of electrons from cold metals. It is the phenomenon that at electric field strengths above some million volts per centimeter electrons are emitted from cold metals. Such high fields occur particularly near sharp points of electrodes. The effect of field emission, therefore, seems to play an important role for the initiation of certain types of spark discharges.

If we discuss field emission by means of the potential well model, we have to superimpose on the potential of the binding forces the

potential of the strong external electric field. The result is the distorted, asymmetrical shape of the potential well edges which is indicated in Fig. 209. Such a potential distribution implies that the potential of an electron on the positive side of the field outside the metal would be smaller than inside the metal. Therefore, electrons can penetrate the potential wall (tunnel effect) at the place indicated by the dotted line and escape into free space. It should be noticed that this explanation of field emission corresponds closely to the explanation of the preionization of atoms (page 186) by a distortion of the potential curve and tunnel effect. Since, according to page 236, the penetration probability increases with decreasing

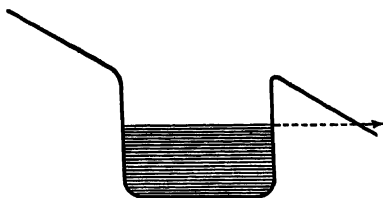


FIG. 209. Distortion of the potential wall in a strong electric field. This makes possible the field emission of electrons from the metal by way of the tunnel effect (dotted arrow).

height and width of the potential wall to be penetrated, one can draw three interesting conclusions. First, the field emission should increase greatly with increasing field strength because, as a result of distorting the edge, the width of the potential wall is reduced. Second, with increasing temperatures the number of escaping electrons should increase because

the electrons in the higher states (bands) in the potential well have lower and thinner potential walls to penetrate (transition to thermionic emission). Third, similar to the case of thermionic emission, the field emission should increase strongly, or begin at lower field strengths, if the effective work function (i.e., the height of the potential wall to be penetrated) is reduced by suitable adsorbed layers on the cathode surface.

The increase of field emission with increasing temperature does not seem to be of particular interest and has not been measured. The increase with increasing field strength has been measured and found to be in agreement with the expectation: the field emission increases exponentially with the field strength. Reliable quantitative measurements are difficult, because adsorbed surface layers and microscopic irregularities cause deviations from the expected behavior. For instance, for thoriated tungsten, an appreciable field emission has been found for fields of 1 million volts/cm, though theoretically field emission was not expected below  $10^7$  volts/cm. However, careful removal of all surface layers and consequent outgassing of the electrode surface led to a decrease of the field emission. The remaining discrepancy seems to be due to the fact that the actual field strength near microscopic peaks of the surface is larger than the computed one. Conclusive proof for this has been recently found by electron-microscope measurements



of emitting metal points. The field emission of cold metals here actually was found to begin at  $3 \times 10^7$  volts/cm.

As a last example for the application of the potential well model we discuss the contact potential between two different metals. Two metals may be brought into contact and then slightly separated. Between them, we can then measure a potential difference, the contact voltage. Each of the metals may be characterized by a potential well with different values of  $\Phi_0$ ,  $\Phi$ , and  $\delta$  (Fig. 210). If the two metals are in contact, they have the same potential. This means that electrons would flow from one to the other, until both electron surfaces have the same height, as it is indicated in Fig. 210. However, now an electron at the surface of metal II has a higher potential energy than one at the surface of metal I, because the work function of metal II in our example is larger than that of metal I. The difference of the potential energies of these electrons is just the potential difference between the two surfaces, i.e., the contact voltage.

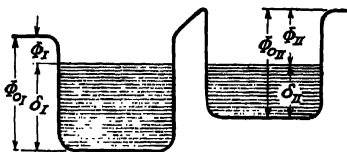


FIG. 210. Explanation of the contact potential between two different metals I and II.

Its value is evidently equal to the difference of the effective work functions of the two metals. This contact voltage exists also between the two separated metals. Because of the difference of the work functions, more electrons will leave the metal I than II, until the two metals have a potential difference which prevents further electron current. Again we arrive at the result that the contact voltage is  $\Phi_I$  minus  $\Phi_{II}$ . This theoretical result is in agreement with experiment, provided that the measurements are made with extremely clean and outgassed metal surfaces. In most cases, surface layers of various kinds cause complications. It is obvious that, conversely, the effect of such surface layers, loose contacts, etc., can be studied by means of analogous considerations with potential well models. This has been done successfully and is of interest for electrical engineering.

### 7-11. Ferromagnetism as a Crystal Property

Ferromagnetism is one of the crystal properties which has been explained by quantum mechanics. That it is actually *a property of the crystal and not of the single atoms* follows from the fact that, for example, iron vapor or iron compounds do not exhibit ferromagnetism (see page 152), whereas the Heusler alloys formed from the nonferrous metals copper, manganese, and aluminum are ferromagnetic. Moreover, certain unmagnetic crystals can be made ferromagnetic just by changing their lattice structures.

The most remarkable property of ferromagnetic materials is that even a very small external magnetic field produces a large magnetization. Moreover, the magnetization does not increase steadily by increasing the field. A large-scale picture of the hysteresis curve (Fig. 211) shows that the magnetization increases in discontinuous steps, the so-called Barkhausen steps. This stepwise, large increase of the magnetization under the influence of a comparatively weak external magnetic field can be explained only by the assumption that *even without an external*

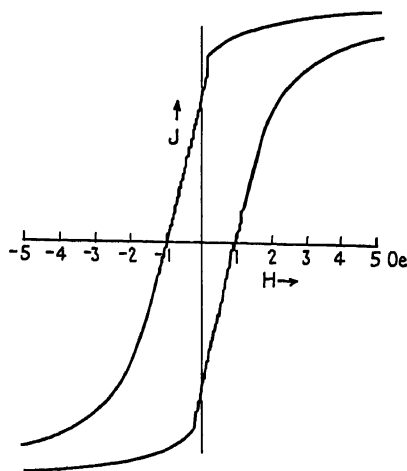


FIG. 211. Magnetization curve with Barkhausen steps. (After Becker and Döring.)

*field entire crystal domains are magnetized (spontaneous magnetization) and that the magnetic field only serves to overcome the resistance which prevents the instantaneous orientation of the magnetic domains in the direction of the external field.*

How can we explain this strong spontaneous magnetization of whole regions in a ferromagnetic substance even without an external field? It can only result from the parallel orientation of the magnetic moments of all or almost all conduction electrons of the highest occupied band in the individual domains in the metal under consideration. Normally, according

to page 439, each state, characterized by the quantum number  $k$ , of the highest energy band is occupied by 2 electrons. According to the Pauli principle (page 169) these two electrons must have opposite spins so that their magnetic moments compensate each other. Now if all magnetic moments are oriented in the same sense, then each  $k$  state can be occupied by only one electron. Then for the same number of electrons twice as many  $k$  states are required as in the nonferromagnetic case. This number of  $k$  states is available in metals, because their highest energy bands are generally only half-filled (page 443). However, the electrons must, on the average, now occupy higher  $k$  states of the conduction band than normally, so that their average zero-point energy is higher too. Now such a spin orientation can occur spontaneously only if energy is *gained* in the process. This means that the energy gain due to the parallel orientation of the spin momenta must overcompensate the increase in the average zero-point energy of the conduction electrons.

The answer to the problem of how a parallel orientation of the spin momenta can occur and lead to a lower potential energy of the system than the normal state of antiparallel spin momenta, was given by Heisenberg in a famous application of quantum mechanics. According to page 231, the energy of interaction of two electrons 1 and 2 which may belong to the ions  $a$  and  $b$ , consists of an electrostatic part which is independent of the spin orientation, and the exchange energy which depends on the mutual spin orientation. It is positive or negative depending on whether the exchange integral

$$A = \int \psi_a(1)\psi_b(2)\psi_a(2)\psi_b(1) \left( \frac{1}{r_{ab}} - \frac{1}{r_{a2}} - \frac{1}{r_{b1}} + \frac{1}{r_{1,2}} \right) d\tau \quad (7-16)$$

is positive or negative. Thus, if this exchange integral of the eigenfunctions of electrons with spin in the same direction is positive, then this ferromagnetic state of the crystal is energetically more favorable than the nonmagnetic state with compensating magnetic moments. From (7-16) it can be seen that  $A$  will be positive if the average distance between the two interacting electrons,  $r_{1,2}$ , is small and the distances between the nuclei and electrons  $r_{a2}$  and  $r_{b1}$  are large. Accordingly we expect, after Slater, a spontaneous orientation of the spin moments of the quasi-free electrons in a crystal, if two conditions are fulfilled:

1. The atoms forming the crystal must not have closed shells, so that the spin momenta of the external electrons are not mutually compensated. Furthermore, these electrons should occupy a state of high orbital quantum number in which the electrons are relatively far from their nuclei.

2. The radius of the electron shell must be small compared to the lattice distance, since only then the integral over the product of the electron eigenfunctions divided by  $r_{1,2}$  has a large value.

Condition 1 is fulfilled for all transition metals in the periodic table, whereas condition 2 holds only for the rare earths and the group, iron, cobalt, and nickel. Though the exchange integral is positive for the rare earths, its value is so small that even a slight thermal agitation is able to counteract the spontaneous orientation of the spin magnets in the crystal. The crystals of the rare-earth elements thus are ferromagnetic only in the neighborhood of absolute zero. Their Curie points, above which their ferromagnetism disappears, is far below room temperature. Only for the metals iron, cobalt, and nickel is the exchange integral large enough to cause a very stable spontaneous orientation of the spin momenta. Only these metals therefore are ferromagnetic up to their fairly high Curie points of 360 to 1000°C. Figure 212 shows the value of the exchange integral  $A$  (7-16) with its depend-

ence on the ratio of the interatomic distance to the radius of the unclosed electron shell. This explanation of ferromagnetism by the atomic theory is confirmed by the fact that ferromagnetism can be produced by increasing the lattice distance of the nonmagnetic metal manganese. According to Fig. 212, its exchange integral is only slightly negative and becomes positive if we increase the interatomic distance by incorporating nitrogen, i.e., by converting the metal to manganese nitride.

A detailed discussion of the behavior of the various modifications (lattice structures) of iron by Dehlinger revealed that apparently not only the exchange energy between an atom and its immediate neighbors

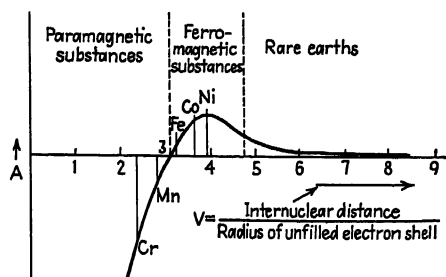


FIG. 212. Curve giving the value of the exchange integral (which determines ferromagnetism and Curie point) for a number of ferromagnetic or nearly ferromagnetic metals as a function of the ratio of interatomic distance to the radius of the incomplete electron shell. (After Becker and Döring.)

is of importance. Because the outer electrons of metals are not well localized (see page 432) or, in other words, because the eigenfunctions are extended over the volume of many atoms, the interaction between a particular atom and its more distant neighbors has to be taken into account and may determine whether, in a particular case, parallel spins are to be expected or not. Consequently, a parallel orientation of the spin momenta of the conduction electrons, and thus ferromagnetism, is to be expected if the exchange integral obtained by summing the exchange integrals between an arbitrary atom and all other atoms of the lattice is positive. With this refinement not only the magnetic behavior of the various modifications of iron can be understood, but also the complicated behavior of the many magnetic alloys which have been discovered in the past 20 years.

The spontaneous magnetization (without an external field) of certain crystals can thus be understood on the basis of the atomic theory. That a piece of iron in general is not magnetic without an external field, is due to the fact that this spontaneous magnetization occurs only in small crystal domains of 100 to 10,000 atoms diameter, and these elementary domains are usually oriented completely at random in the

crystal. An external magnetic field is thus required in order to orientate the domains. That such an orientation is possible means that these elementary domains can be rotated in the solid crystal though they offer some resistance against an arbitrary rotation. This peculiar crystal property has been studied carefully by measuring the hysteresis curves of single crystals of iron. For these studies the external magnetic field was directed along an edge, along a surface diagonal or the volume diagonal of the cubic iron crystal. Figure 213 shows the results.

If the magnetic field acted along a cube edge (100), the magnetization increased rapidly with increasing orienting field to the saturation value  $M_s$ . Evidently only a very small orienting force is sufficient to turn the spontaneously magnetized domains in the direction of the cube edge. If the direction of the field coincided with a surface diagonal (110), the magnetization increased rapidly up to the value  $M_s/\sqrt{2}$ , and with the field in the direction of the cube diagonal (111), up to the value  $M_s/\sqrt{3}$ . It then rose slowly to the saturation value  $M_s$ . This behavior is understandable if it is assumed

that all spontaneously magnetized domains do not offer any appreciable resistance against being rotated in the direction of the cube edge closest to the field direction, but that it requires more force to turn the spins from this preferred direction into the actual field direction. After the external field has been removed, the domains will reorient themselves in the direction of the nearest cube edge and will remain there if the thermal agitation is not too great. This is the explanation of the well-known remanent magnetism.

For the usual polycrystalline iron, and particularly for iron as it is treated in many manufacturing processes, the behavior is much more complicated than for single crystals. This is chiefly because the inner stresses produced in working the iron give rise to preferred directions of orientation for the domains. However, even here the complicated phenomena of ferromagnetism can be understood in principle from the well-known rules of atomic physics.

Recently the interesting question has been studied how large a Fe microcrystal must be in order to exhibit ferromagnetism (in contrast

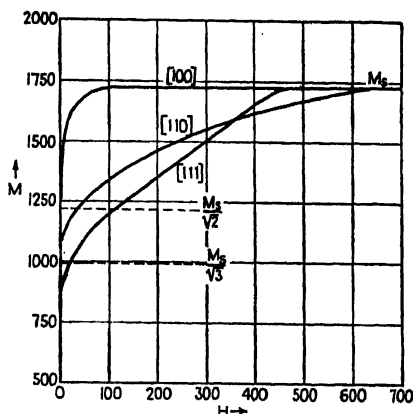


FIG. 213. Magnetization curves of an iron single crystal for orientation of the field in different crystal directions. (After Honda and Kaya.)

to the Fe atom). According to the experiments of König the micro-crystal must have at least 64 elementary cells (page 425). If the surface atoms which are not essential for the spin orientation are disregarded, this number reduces to 8, i.e., to a body which also seems to play a special role as a crystal seed. Here we have apparently arrived at the borderline between atomic complexes and solid bodies, i.e., micro-crystals.

One more important property of metals should be mentioned at this point, though it cannot yet be explained by atomic physics: *supraconductivity* (page 445). This strange effect also depends on a special

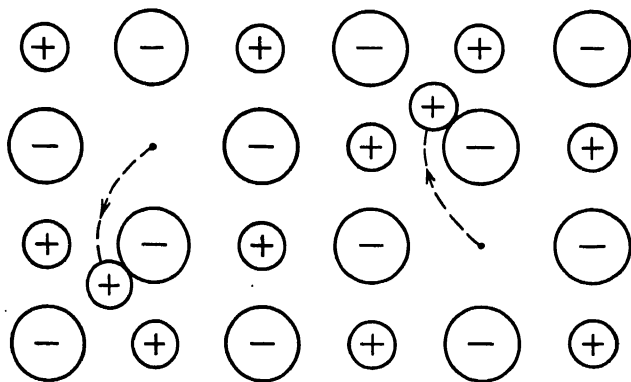


Fig. 214. Positive interstitial ions and vacant lattice places in an ionic crystal: Frenkel defects.

electron orientation and apparently, just as ferromagnetism, is possible only in certain crystal lattices. It seems that, because of a special "configuration," the conduction electrons can move through the supraconducting metallic lattice practically without any resistance. Since there is no satisfactory atomic explanation for this important and highly interesting phenomenon, we have not devoted a special section to it, but only mention it in concluding the discussion of ferromagnetism which also depends on a special electron orientation in certain crystal lattices.

### 7-12. Lattice Defects and Their Bearing on Diffusion and Ionic Conduction in Crystals

Ferromagnetism and supraconductivity seem to be structure-sensitive properties (see page 422) in so far as a particular lattice is necessary for these phenomena. We shall now discuss those properties of solids which, like diffusion and ionic conductivity, depend essentially on lattice defects. In an ideal crystal each constituent is in its own poten-

tial trough and any migration of a particle—for example, the exchange in the location of two particles (atoms, ions, or molecules according to the lattice type, page 423) in the case of spontaneous diffusion—requires energy in order to overcome the potential wall. Spontaneous diffusion therefore is strongly temperature dependent. Because it increases exponentially with the temperature, it is difficult to observe at room temperature.

According to page 422, each real crystal has a large number of lattice

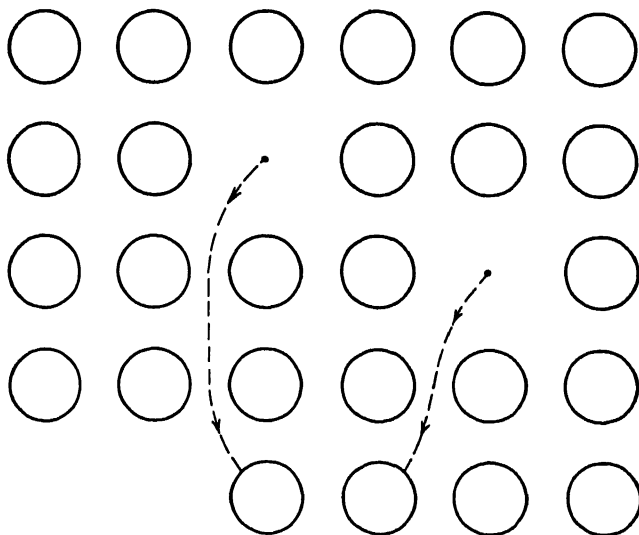


FIG. 215. Positive and negative vacant lattice places: Schottky defects.

defects, such as vacant lattice places which would normally be occupied by a constituent, and the so-called interstitial ions. During the growth of the crystal, according to Fig. 214, a normal lattice ion can “erroneously” be “frozen” in an interstitial place, i.e., it did not get into its right potential trough but, by deformation of the crystal, has made a place for itself in the neighborhood of the correct place. By this process there is produced in the crystal an empty place and an interstitial ion. This type of defect is called a *Frenkel defect*. Also during crystal growth or by subsequent migration, ions of both signs (since the crystal is electrically neutral) may have come to the crystal surface; empty places then remain in the crystal (Fig. 215). These defects are called *Schottky defects*. Both types of defects enable lattice constituents to have a much greater mobility, since *with lattice holes the diffusion can progress stepwise through the lattice holes*. This is called *hole diffusion*, and the process is called *defect conduction*. Obviously much less activa-

tion energy is required for hole migration or the motion of interstitial ions than for an exchange of normal ions in a lattice. Since the number of lattice defects depends on the incidentally more or less regular crystal growth and its later treatment (heating, mechanical stress, etc.), the defect diffusion varies greatly between different samples of the same type of crystal. This behavior is in contrast to the characteristic high temperature diffusion mentioned above, which is independent of lattice defects.

The study of the diffusion of atoms or ions in their own crystal, for instance of lead atoms in lead metal, has become possible, as mentioned on page 309, by the radioactive tracer method (diffusion of radioactive atoms or ions). The diffusion of two solid materials into each other, when in contact, e.g., that of silver into gold and conversely, is one of the migration phenomena which is easier to study. At low temperatures the diffusion in solids is exclusively due to vacant lattice places and interstitial atoms. However, it is obvious that this migration process is also strongly temperature-dependent. With increasing temperature higher and higher lattice vibrations are excited and lead to a "loosening up" of the lattice which facilitates the migration of lattice constituents. It is, finally, understandable that the velocity of diffusion of one solid into another is the larger, the less the difference in the diameters of the two kinds of atoms. An occasionally very high concentration of lattice defects plays an essential role in the semiconductors and phosphors which we shall discuss in detail later. In these crystals, either one of the lattice constituents is present in excess (e.g., Zn in ZnO), or foreign atoms are introduced into the lattice (so-called impurity-activated crystals). In both cases we find interstitial atoms as well as lattice vacancies.

It is evident that gas atoms or ions can also diffuse into solids. The diffusion of hydrogen through palladium and iron, to cite only two prominent examples, is well known to gas discharge physicists.

The great importance of diffusion processes for the practical use of metals becomes apparent from the example of rusting of iron. As soon as metallic iron is covered by a layer of rust, it is no longer possible for oxygen or OH radicals to attack the metal since this is now protected by a layer of another solid, rust. Further rusting is possible only by diffusion of either oxygen or metal through the crystalline layer of rust. Because of the smaller radii of the positive ions it seems that the diffusion of the iron ions to the surface of the rust layer is more important than the reverse diffusion of oxygen. The number of similar examples can be increased arbitrarily.

We have not taken into consideration up to now that the lattice constituents of an ionic lattice are ions of opposite polarity. Thus, if



we apply an electric field to an ionic crystal, its ions migrate to the electrodes of opposite sign, so that ionic crystals exhibit an electric conductivity which is ionic just as is the conductivity in electrolytic solutions. It is obvious that this ionic conductivity has nothing to do with the migration of electrons and the corresponding electronic conductivity in certain crystals (semiconductors) which we shall discuss later. Because the ionic conductivity depends on the migration of lattice ions, we have the empirically well known negative temperature coefficient of resistance of ionic crystals: the resistance decreases with increasing temperature, because the lattice loosens up. *At low temperatures the ionic conductivity is strongly structure-dependent, and the migration occurs almost exclusively via the vacancies and interstitial lattice places. At high temperatures, on the other hand, the structure dependency vanishes and, because of the loosening of the lattice, the ions migrate by exchanging places in the lattice.*

That the ionic conductivity increases exponentially with the temperature is to be expected for an exchange of places by overcoming potential walls. For a particular sample of a NaCl crystal, according to Smekal, the conductivity could be represented by the formula

$$\sigma = 0.1e^{-10,000/T} + 10^6e^{-24,000/T} \quad (7-17)$$

The difference in the order of magnitude of the constants, 0.1 and  $10^6$ , shows that only the  $10^7$  part of the lattice constituents participates in the conductivity represented by the first term. This term consequently describes the conductivity due to lattice defects. The second term represents the conductivity resulting from the migration of normal lattice ions. In agreement with this explanation the activation energy of the defect conduction is only 40 per cent of that of the second term (10,000 to 24,000). The average height of the potential walls for both processes is 0.8 and 2.0 ev, respectively. Because of the small height of the potential wall (activation energy) in the exponent, the first term of the conductivity formula predominates at low temperatures in spite of the small numerical factor, whereas at high temperatures the second term, because of its large factor, predominates. Since, in the general case, holes, interstitial ions, and normal ions of both signs contribute to the conductivity, the temperature dependence of the conductivity of ionic crystals should be represented by a formula with four terms. Practically, however, it can be represented by two terms. Because of the different activation energies of the various partners and the limited temperature regions in which it is feasible to make measurements, only two terms play an important role, unless the accuracy of the measurements is increased considerably.

### 7-13. Light Absorption and Electron Motion in Halide Crystals. The Primary Photographic Process

In the last section we were concerned with the movement of lattice constituents as well as of foreign atoms or ions in the crystal lattice and learned of the decisive influence of lattice defects on this migration. We shall now discuss the effect of light on ionic crystals, and here we shall deal for the first time with the *migration of electrons in nonmetallic crystals*. The fundamental investigations on the effect of light on ionic crystals were carried out by Pohl, Hilsch, and coworkers on "colored"

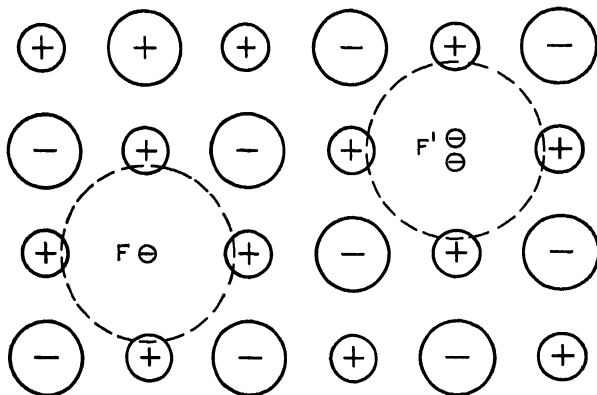


FIG. 216. Schematic representation of the probable nature of  $F$  and  $F'$  centers in colored alkali-halide crystals: one or two electrons in the lattice place of a missing negative ion.

alkali-halide crystals. In addition to their general interest, they are important for understanding the fundamental processes of photography.

If transparent alkali-halide crystals are irradiated with sufficiently short ultraviolet light, with X- or cathode rays, or if they are heated in the vapors of their characteristic metals (e.g., NaCl in Na vapor), the crystals become colored. The characteristic color for each crystal varies, for instance, from yellow for LiCl to blue for CsCl. The absorption band corresponding to this color, which is missing in the untreated crystal, has been called the  $F$  band (Farb band or color band) by Pohl. The corresponding absorbing centers are called  $F$  centers. It required many years of intensive studies to explain the nature of these  $F$  centers, and it is not possible to present here all the experiments which support the presently accepted explanation. According to our present knowledge, an  $F$  center is not, as was formerly believed, simply a metal atom produced by attaching a free electron to an arbitrary metal ion of the lattice. An  $F$  center seems to consist of an electron which has taken the place of a missing negative halogen ion (Fig. 216) and then, as a non-

localized electron (page 415), alternately forms an absorbing Na atom with each of the six surrounding Na ions. Its excitation energy of about 3 volts, as computed by wave mechanics, agrees well with the wavelength of its absorption band. To ionize an  $F$  center, i.e., to release the electron from its potential well, an energy about 0.1 volt higher than its excitation energy is required. The following experiment is in agreement with this result: If a colored crystal at a sufficiently low temperature is irradiated with the wavelength of its absorption band, then the  $F$  electron excited by the absorption remains bound to its place and the crystal thus remains an insulator. On the other hand, at elevated temperatures the energy lacking for the complete release of the  $F$  electron is supplied by the thermal vibrations of the lattice and the  $F$  electron is raised to the previously unoccupied conduction band. Here, as a free electron, it is responsible for electronic conductivity if an electric field is applied to the crystal (Fig. 217); the crystal exhibits "photoelectric conductivity."

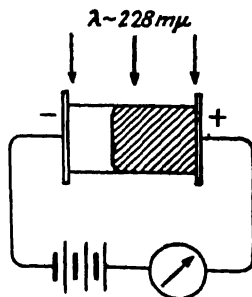


FIG. 217. Diagram of the Pohl experiment for the electric measurement of the electron migration in a colored  $\text{KHI-KBr}$  mixed crystal. By ultraviolet irradiation electrons are liberated and migrate in the electric field from one vacant negative hole to the other toward the anode, thus establishing an electric convection current.

Pohl furthermore found that the destruction of  $F$  centers by irradiating in the  $F$  band at not too low temperatures produced not only the expected reduction of the absorption intensity of the  $F$  band, but also the appearance of a new longer wave absorption band which he called  $F'$  band. The nature of these  $F'$  centers which are responsible for this longwave absorption has been explained in the following way: the  $F'$  center is supposed to be an  $F$  center that has captured an additional electron (Fig. 216), or, in other words, two instead of one electron in the place of one missing Cl ion. The absorption of the  $F'$  bands lies at longer wavelengths than that of the  $F$  bands, because it requires only a small amount of energy to separate the loosely bound second electron. The formation of  $F'$  centers according to the present explanation seems clear: an  $F$  electron, released from its potential trough by the absorption of light with the cooperation of the thermal lattice energy, is captured by another  $F$  center which is then transformed into an  $F'$  center. In agreement with this explanation, two  $F$  centers are destroyed for each quantum absorbed in an  $F$  band (quantum yield 2). One is destroyed directly by absorption and one is transformed into an  $F'$  center. The fact that the mean free paths of the photoelectrons released from the  $F$  centers is inversely proportional to the density of  $F$  centers is

further direct confirmation of the explanation of the  $F$  centers as presented here. Furthermore, in general it is not possible to produce more than  $10^{18}$   $F$  centers per cubic centimeter in a crystal. This is further evidence in favor of the present theory, because in general it is not possible to form more  $F$  centers than there are  $\text{Cl}^-$  vacancies. Consequently, *the production of  $F$  centers consists in the introduction of free electrons into the halogen vacancies in the lattice.* The free electrons are produced by the absorption of ultraviolet light or X-rays or by collisions



FIG. 218. Pohl's experiment by which the migration of electrons in an alkali-halide crystal is made visible. At  $600^\circ\text{C}$  the electrons in a KBr crystal migrate so quickly toward the positive pole that the rear boundary of the "blue cloud" formed by the  $F$  centers moves visibly toward the anode. (After Pohl.)

of cathode-ray electrons. If the coloration is produced by heating the crystal in metal vapor, the metal ions attach themselves to the crystal surface (together with halogen ions which migrate outward and thus produce vacant lattice places), while only their electrons migrate into the crystal. That actually an electron migration into the crystal produces the  $F$  centers, and thus the coloration, was shown by Pohl in his famous electrical experiment. If a potential difference of several hundred volts is put across a completely transparent KBr crystal and the lattice "loosened" by raising the temperature to 400 to  $600^\circ\text{C}$ , then electrons migrate from the negative electrode into the crystal and fall into  $\text{Br}^-$  vacancies, thus producing absorbing  $F$  centers. A blue cloud can be seen moving into the crystal, which can be moved back and forth as the polarity of the electric field is reversed (Fig. 218). Pohl called this produc-

tion of  $F$  centers by introducing electrons into the crystal "*additive coloration*" in contrast to the coloration of crystals by irradiating them with X-rays. In this last case the electrons can return to their original places when the crystal is heated or irradiated with visible light. *The crystal can be bleached, in contrast to an additively colored crystal.* All these experiments with alkali-halide crystals have not only revealed a number of interesting processes in solids but they have also clearly shown *that lattice defects are of decisive importance for the effects of light on ionic crystals and the behavior of electrons in such crystals.*

The same situation holds for the basic atomic processes of photography which are so complicated because of this structure dependency. Free electrons are primarily produced by absorption of light in the AgBr microcrystals of the photographic emulsion just as in the alkali-halide crystals. After filling the lattice vacancies, these electrons attach themselves to  $\text{Ag}^+$  ions and form absorbing Ag atoms. The decisive but not yet entirely understood process is now that the individual Ag

atoms coagulate and form colloidal-like silver nuclei. After the development of the "latent" image, i.e., after the chemical reduction of further Ag atoms, these silver nuclei cause a blackening in the plate at those places which had been irradiated.

Theoretically, in producing the latent image, each absorbed photon can release one electron which attaches itself to an ion and forms a silver atom. This quantum yield of 1 can actually be obtained in the elementary photographic process according to measurements of Eggert and Noddack. In general, the quantum yield depends on the illumination during the exposure since too strong exposures produce a reformation of AgBr which reduces the effective quantum yield.

That the addition of easily ionizable molecules, such as  $\text{Ag}_2\text{S}$ , increases the sensitivity of the photographic emulsion can be understood without further discussion. It is also clear that a certain minimum energy is necessary to separate the electron from the halogen ion. This energy determines the wavelength limit beyond which the emulsion is no longer sensitive. Photographic emulsions are always sensitive to shortwave light but, in general, they are not sensitive to longwave red light. It is known, however, that photographic emulsions can be sensitized for longwave light by adding certain dyes which are able to absorb longwave light.

Although one has been able to explain to a certain extent the elementary photographic process, one cannot yet fully explain the development of the latent image. It is known that the number of primary Ag atoms which form the latent image is increased by a factor up to  $10^8$  by chemical reduction of silver atoms, and that the existence of silver nuclei is necessary for this chemical development. Yet nothing definite is known about their structure or effects. In particular, it is not yet known whether the nucleus consists of colloidal silver or, as many investigators believe, regularly built micelles in which foreign atoms play an important role. The atomic physicist is still looking forward to more complete empirical results.

#### 7-14. Electronic Conduction in Semiconductors and Related Phenomena. Barrier Layers and Rectification, Crystal Photoeffect. The Transistor

In the last section we became acquainted with the electronic conductivity of insulating crystals which depends upon the motion of free electrons. These were produced, i.e., raised to the normally unoccupied conduction band, by the absorption of light in the crystal. *In some solids electrons can be raised to the conduction band not only by the absorption of light but also by the thermal vibrations of the lattice, if the temper-*

ature is not too low. This gives rise to a "dark conductivity." These materials are called *semiconductors*.

By definition a semiconductor is a crystal which is an insulator at absolute zero but which, at higher temperatures, has an electronic conductivity that increases exponentially with the temperature. In order to determine whether the conductivity is the result of the migration of electrons and not of ionic diffusion, the *Hall effect* is used. In Fig. 219 an electric field is applied in the  $y$ -direction of a semiconductor

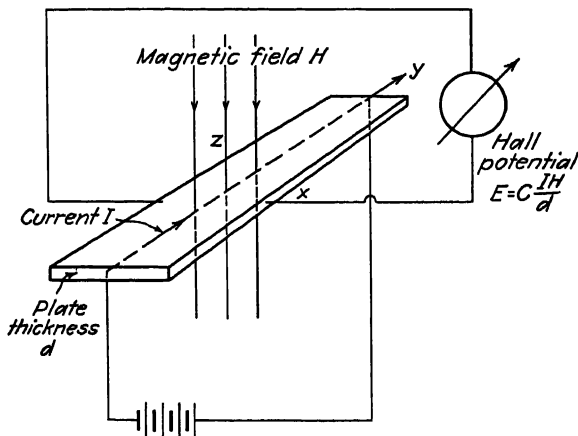


FIG. 219. Schematic sketch of the arrangement for measuring the Hall voltage of semiconductors.

slab lying in the  $xy$ -plane, and a magnetic field acts in the  $z$ -direction. Because of their mobility, the electrons moving along the  $y$ -direction are then deviated by the magnetic field in the  $x$ -direction. The result is a potential difference, the so-called *Hall potential*, perpendicular to the direction of the current. This potential difference can be observed only in the case of electronic conductivity, but not for conduction by ions, because, according to Sec. 7-12, the ion mobility is too small. *The Hall effect thus can be used to distinguish between semiconductors and ionic conductors.*

*Electronic conductivity in a nonmetallic solid (semiconductor) thus is due to electrons which are released by thermal lattice vibrations (similar to thermal collisions) and raised to the conduction band which is unoccupied at  $T = 0$ . This is possible in the case of a pure crystal only if the energy difference between the normally empty and the normally filled energy bands is not too large compared with the thermal energy  $kT$ , i.e., if it is at most several tenths of a volt in comparison with the 2 to 10 volts in most crystals. Such a crystal with a small gap between the two bands*

(e.g., germanium with a gap of 0.75 volt) is called an *intrinsic semiconductor*. A crystal which is not an intrinsic semiconductor can be made into a semiconductor (and this is by far the most frequent case) by introducing into it foreign atoms, or a stoichiometric excess of one of the component atoms which has a sufficiently small ionization energy in the crystal. In the energy band model this means that, according to Fig. 220, the energy states *D* of the valence electrons of these interstitial atoms are situated so closely below the normally unoccupied conduction band *B*, that electrons can get from *D* into the band *B* as a result of the thermal ionization of the metal atoms. This type of crystal is called

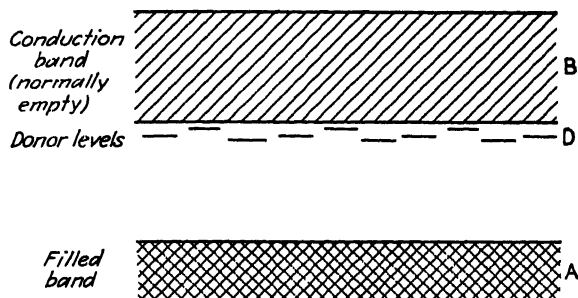


FIG. 220. The energy band diagram for an *N*-type semiconductor: *A* = filled energy band; *B* = empty conduction band; and *D* = impurity levels of "donor atoms," from which electrons can easily reach the conduction band *B*.

an *excess semiconductor*, because it requires an excess of metallic ions in the lattice to provide the electrons, or an *N-type semiconductor* because negative charges are released from the *D* levels. The *D* levels occasionally are called *donor levels*, because they donate the conduction electrons. The third possible type of semiconductor is called the *P-type or defect semiconductor* because it has a defect of metallic atoms in the lattice or an excess of the electronegative constituents (e.g., oxygen in the metal oxide semiconductors). The energy states of these electronegative atoms lie closely above the band *A* of the crystal which is completely occupied by electrons (Fig. 221). Only a small amount of energy is necessary to capture an electron, i.e., to bring it from the crystal energy band *A* to the normally unoccupied levels *C* of the electronegative atoms, where it is held fast. The "holes" which are thus produced in the previously full band of the crystal have an effective positive charge and are mobile. When an electric field is applied, a current flows which apparently is conducted by positive charges (positive holes, therefore *P-type semiconductor*). The electronegative atoms in the *C* levels are often called *receptor atoms*, because they receive the electrons from the filled band and thus cause the positive holes.

The excess or defect semiconductors are also called reduction or oxidation semiconductors, respectively, because in the practically important metal oxide semiconductors the excess of the metal or oxygen atoms is produced by chemical reduction or oxidation. Consequently, depending on the preliminary treatment, the same crystal can be an excess or a defect semiconductor, a typical indication for the "structure dependency" of the properties of a semiconductor. Examples of

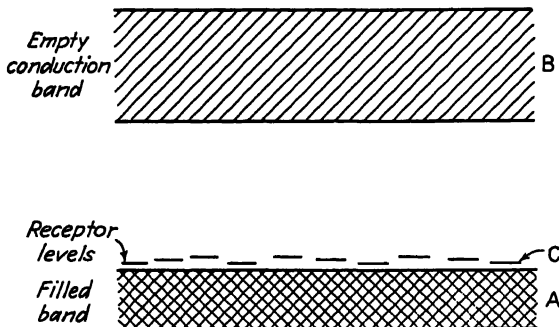


FIG. 221. The energy-band diagram for a *P*-type semiconductor: *A* = filled energy band; *B* = empty conduction band; and *C* = empty impurity levels of "receptor atoms," which can easily be filled from *A*, thus producing mobile positive holes in *A*.

intrinsic semiconductors are pure crystals of silicon and germanium which have recently been extensively investigated. The most important examples of practical excess and defect semiconductors are, respectively, zinc oxide,  $\text{ZnO}$ , which is heated in vacuum, and cuprous oxide,  $\text{Cu}_2\text{O}$ , heated in oxygen.

Since the energy required for the transition  $D \rightarrow B$  (*N* type) or  $A \rightarrow C$  (*P* type) in the so-called dark conductivity of semiconductors is supplied by thermal collisions, the conductivity can be represented by a formula of the general form

$$\sigma = Ae^{-W/kT} \quad (7-18)$$

in which  $W$  is the "activation energy" of the process (the energy difference from  $D$  to  $B$  or  $A$  to  $C$ ) and is in general, just as is the constant  $A$ , a function of the temperature, since  $W$  depends on the loosening of the lattice which is determined by the temperature.

On the basis of this theory of semiconductivity, a large number of phenomena of physics, many of which have been known for a very long time, can now be explained. For instance, for many decades, *barium oxide cathodes* have been used in many electronic instruments, because it had been found that at a comparatively low temperature barium oxide has a much higher thermionic electron emission than any metal.



This is also a semiconductor effect. By the "formation" of the barium oxide cathode, an excess of metal atoms is produced in the microcrystals, and these excess metal atoms, being in the  $D$  levels of Fig. 220, release electrons into the conduction band  $B$  at a sufficiently high temperature. Now the energy necessary to release these electrons from the conductivity band  $B$  into free space (effective work function of the semiconductor) is much smaller than that for their release from the normally occupied bands of a metal. This reduced effective work function appears in the exponent in Richardson's equation (7-15) and thus causes, at the

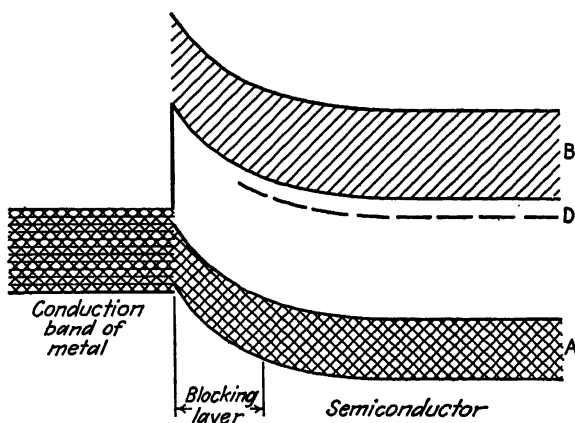


FIG. 222. Energy-band diagram for an  $N$ -type semiconductor in contact with a metal; formation of a barrier layer with rectifier properties.

same temperature, an increase of the electron emission by many orders of magnitude, as compared with that of a metal.

The well known *rectifier and detector properties of numerous metal oxide crystals are also semiconductor phenomena* and can be explained with the aid of the energy band diagram Fig. 220 or 221. If an  $N$ -type semiconductor is brought into contact with a metal electrode, electrons from the high impurity levels  $D$  "flow down" into the metal, until the electric field set up in the boundary layer causes such a distortion of the energy bands (Fig. 222) that further electron current is impossible, at least without an external field. However, if an alternating electric field is now applied across this metal-semiconductor contact, the energy bands on the right side of Fig. 222 are raised and lowered, alternatingly, with respect to the energy band of the metal electrode at the left. Now with positive polarity of the metal electrode, the bands of the  $N$ -type semiconductor are raised to such an extent that electrons from the  $D$  levels can again flow down into the metal. When the metal electrode is negative, this is impossible, unless the potential applied to

this "blocking layer" becomes too high. In this case the energy bands of the semiconductor at the right side of Fig. 222 are lowered so much that electrons from the metal can penetrate the blocking layer by the tunnel effect (page 234) and flow into the  $B$  levels of the semiconductor. Thus this theory requires an easy electron flow from the semiconductor to the metal and a blocking of the current for flow in the opposite direction. However, the resistance of the blocking layer with this polarity should decrease with increasing applied potential. This behavior has been found empirically and holds for most semiconductor rectifiers with the exception that some of them (like cuprous oxide,  $\text{Cu}_2\text{O}$ ) are  $P$ -type semiconductors, so that we have to use Fig. 221 for our discussion. The detector property of semiconducting crystals is, of course, coupled with rectification and thus need not be treated separately.

According to our above treatment, the decisive phenomenon of  $N$ -type semiconductivity is the thermal release of electrons from the impurity levels  $D$ . However, the transitions  $D \rightarrow B$  or  $A \rightarrow C$ , which lead to the production of free electrons or free positive holes, respectively, in the semiconductor lattice, need not necessarily be caused by thermal collisions (lattice vibrations), but they can also be caused by light absorption. Consequently, all these semiconductors show weak, longwave absorption bands which can be attributed to these transitions. Absorption of the corresponding radiation by the alkali halides (page 459) produces an electric conductivity which, in contrast to dark conductivity, is called photoconductivity. This release of electrons in the interior of semiconductors (and by shortwave light in other crystals also) is called the "internal photoelectric effect."

Two different methods have been developed to make use of this release of photoelectrons in semiconductors for the detection of radiation. This radiation might cover the whole range from the infrared ( $3$  to  $6\ \mu$  with  $\text{PbS}$ ,  $\text{PbSe}$ , and  $\text{PbTe}$  semiconductors) down to the  $\gamma$ -rays. Either the photoelectric conductivity of the semiconductor is measured directly in a bridge as a function of the incident radiation, or semiconductor photocells are used which themselves produce a photocurrent that is proportional to the number of incident photons. If, in different layers of a semiconducting crystal, very different numbers of electrons can be released photoelectrically, then we get a concentration gradient of free electrons in the interior of the crystal. The crystal acts as an electron-concentration element with an emf which, because of its origin, is called a *photo-emf*. In  $\text{Cu}_2\text{O}$  the photo-emf can amount to as much as  $0.4$  volt. If a semiconductor is placed between two metal contacts and these are connected to a microammeter, a photocurrent flows when

the crystal is irradiated by light: we then have a *semiconductor photoelement*. These have become important for lux meters and other technical applications. The necessary electron concentration gradient can be obtained in two different ways. First, varying concentrations of lattice defects in different layers of the semiconductor can be produced by proper heat treatment, so that varying amounts of electrons are released when the crystal is irradiated. Second, a crystal can be selected which has layers so thick that the incident light will be strongly absorbed. Then many electrons will be released close to the upper surface where the illumination is strong, whereas on the lower side the number of released electrons will be small because only a small portion of the light reaches it. Technically, the first type of photoelement is preferred.

The basic processes of this phenomenon have been carefully studied, in particular for cuprous oxide (a *P*-type semiconductor) where the necessary *C*' levels (Fig. 221) are due to excess oxygen atoms in the  $(Cu_2O)$  lattice. This is occasionally described as a partial chemical transformation of  $(Cu_2O)$  to  $(CuO)$  by tempering in an oxygen atmosphere. *By this tempering the photoelectric conductivity of pure  $(Cu_2O)$  can easily be increased by eight powers of ten!* If a thin layer with only a few excess atoms and thus a high insulation (*barrier layer* or *blocking layer*) is formed on a crystal layer of good conductivity, and if over these a thin transparent metallic electrode is placed, we have a *semiconductor photoelement* or *barrier layer photoelement* (Fig. 223). A similar widely used photoelement has been developed using selenium, but little seems to be known about its basic atomic processes.

We have briefly discussed above the origin of the internal photo-emf. Actually the theoretical problem of the photoelements is much more complicated, because the processes in the barrier layers, which are formed at the contacts between metal and semiconductor, have to be taken into consideration. We have pointed out some of these barrier layer problems above in our discussion of the rectification properties of semiconductors.

We mention finally one of the most recent developments in the field of semiconductivity which combines scientific interest with practical importance. It is the so-called *transistor*, developed in the Bell Tele-

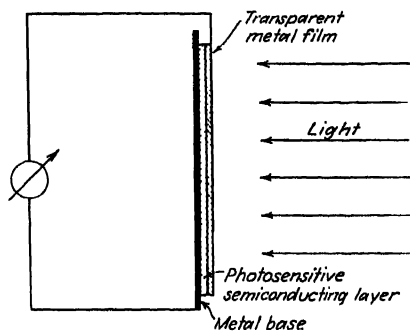


FIG. 223. Diagram of a barrier layer photocell with electric circuit.

phone Laboratories by Bardeen and Brattain. The transistor is essentially a germanium crystal diode for current amplification. Figure 224 shows the schematic diagram. A small germanium crystal rests on a metal base, whereas two point contacts touch its surface at a distance of 0.05 to 0.25 mm from each other. A current of the order of 1 ma flows in the emitter circuit. As a surface layer of  $10^{-5}$  cm thickness of the crystal is a *P*-type semiconductor, positive holes are emitted from the positive point contact. A part of this hole current, which is modulated with the frequency of the signal to be amplified, reaches the negative collector point and changes its current-voltage characteristic. It is known

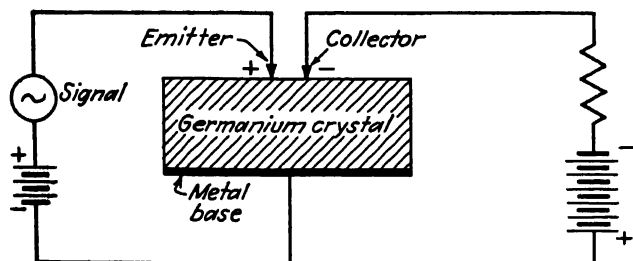


FIG. 224. Electric circuit for the transistor. (After Bardeen and Brattain.)

that in the surface layer of the germanium a barrier layer exists which prevents the flow of electrons from the metal contact into the crystal. The space charge in the barrier layer under the collector point now is changed by the positive holes from the emitter in such a way that the collector current is increased or reduced with the signal frequency.

In addition to its probable future role in electronics, the transistor is of interest because its theoretical explanation required a thorough study of the semiconductor properties of germanium which, depending on its treatment, its temperature, and the potential of its metal contacts, shows conductivity due to electrons as well as holes. The electrons and holes either come from the *C* and *D* levels, or from transitions between the filled band *A* or the conduction band *B*, respectively, and the conduction band of the metal in contact with the semiconductor. For further details the reader is referred to the literature. However, it is clear from this brief discussion that physics of semiconductors has become one of the most important fields of solid state physics.

### 7-15. Luminescence of Crystal Phosphors

The last important example of processes which depend essentially on the microstructure in the crystal lattice is crystal luminescence, which is closely related to semiconductivity. Many crystals show lumines-

cence after irradiation with shortwave light,  $\alpha$ -,  $\beta$ -, or  $\gamma$ -rays, if they have been "activated" by addition of a small percentage of metal atoms. The term luminescence here comprises the phenomena of *fluorescence* and *phosphorescence*. Fluorescence is the phenomenon of light emission immediately after radiation has been absorbed in the crystal, i.e., after the normal period of  $10^{-8}$  sec. Phosphorescence, on the other hand, is the more complex process by which light of the same or longer wavelength is emitted for a comparatively long time after irradiation. From these definitions it is apparent that fluorescence and phosphorescence are not always easy to distinguish. Physically, however, *phosphorescence involves a storage of the absorbed energy before its reemission*. Luminescence can be observed in large single crystals: mostly, however, a microcrystalline powder is used.

From what has been said so far, *ideal crystals are not luminescent*. However, certain crystals do show luminescence, particularly at low temperatures, in spite of the fact that they apparently consist of chemically pure homogeneous materials. These crystallize simultaneously in different types of lattices, so that they contain lattice defects and internal surfaces which are essential for phosphorescence. *All other crystals with phosphorescence at normal temperatures consist of a lattice with interstitial or substitutional impurity atoms which serve as emitting centers*. The best known crystal phosphors are zinc and cadmium sulfide,  $\text{ZnS}$  and  $\text{CdS}$ , which are used as single or mixed crystals, especially if they are activated by copper or manganese. There is approximately one emitting metal atom for every  $10^4$  lattice atoms. The oldest phosphors which were investigated carefully by Lenard are the alkaline-earth sulfides and oxides activated by  $\text{Cu}$ ,  $\text{Mn}$ ,  $\text{Pb}$ , or rare earths. Certain silicates such as  $\text{Zn}_2\text{SiO}_4$  activated by  $\text{Mn}$ , fluorspar,  $\text{CaF}_2$ , activated by heavy metals or rare earths, and numerous tungstates activated by  $\text{Cr}$  and  $\text{Mn}$  are also good phosphors. Of especial scientific interest are the alkali halides with heavy metals and rare earths which were investigated by Pohl and Hilsch, because, in contrast to the above-mentioned phosphors, they can be grown in large crystals and therefore can be investigated under more perfect conditions.

The fundamental process of luminescence can be described in a few words. *By absorption of light an electron leaves the highest occupied energy band of the basic lattice and enters a higher unoccupied band (conduction band). It then moves around in the lattice by thermal diffusion until after a shorter or longer time it comes to an activator atom or ion and gives up its energy by emitting light*. In contrast to fluorescence, however, *in phosphorescence we always have a temporary storing of the excitation energy by trapping of the electron in some vacant lattice position or*

*similar electron trap.* The dependence of phosphorescence on the structure is obvious from the role of the trapping centers. The temperature dependence is also easy to understand. At low temperatures the trapped electrons have difficulty in getting out of the traps again, because the thermal energy of the lattice vibrations is so small. Only a few electrons are thus released per second, and we have a persistent, but weak, emission which decays exponentially. On the other hand, at higher temperatures the trapped electrons are released at a higher rate from their traps and thus produce a more intensive, but rapidly decaying, phos-

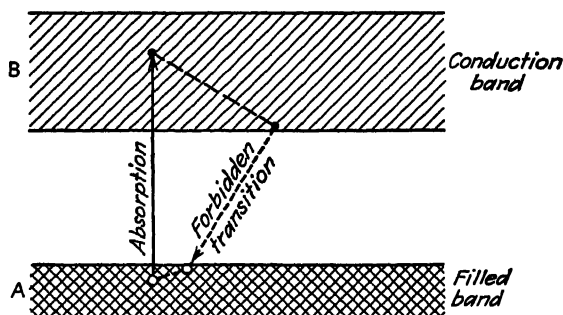


FIG. 225. Energy-band diagram indicating why pure ideal crystals generally do not exhibit any luminescence.

phorescence. The light sum (the total light emitted during the decay) for the same excitation is the same in both cases, as is to be expected.

The main features of the great, almost confusing, number of experimental results which experimenters have assembled since the time of Lenard now can be explained. First we ask why an ideal lattice does not exhibit luminescence. This can be explained by means of Fig. 225. If an electron of the occupied energy band *A* absorbs a photon, it is raised to that state of the higher unoccupied band *B* which has the same *k*-quantum number. Now if it does not immediately jump back under emission of light (fluorescence), it will lose energy by interacting with the lattice, until it reaches the lower edge of the energy band *B*, where it has a different quantum number *k*. This transfer of energy occurs very rapidly because the period of a lattice vibration with  $10^{-14}$  sec is very short compared with the natural life of  $10^{-8}$  sec in the upper state. The direct reverse transition mentioned above, therefore, is very improbable. Now, however, the electron *cannot* jump back from the lower edge of the upper band into the band *A* by emitting a photon, because the corresponding quantum state *k*, to which a transition could occur, is occupied. Luminescence by ideal, pure crystals therefore seems impossible.

The luminescence of metal-activated crystals can be explained by means of the energy band diagram Fig. 226 which is closely related to that of a crystal with *N*- and *P*-type semiconductivity. *A* is again the highest occupied energy band and *B* the first normally unoccupied conduction band of the normal crystal lattice. In addition, there are again "impurity levels" *C* and *D* of substitutional or interstitial foreign atoms. *C* are the activator states which, in general, are occupied by electrons, because transitions to the completely filled band are impossible. The states *D*, finally, belong to electron traps, in which elec-

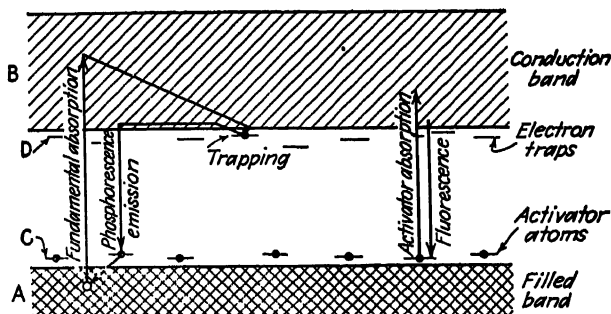


Fig. 226. Energy-band diagram of a crystal with activator atoms and electron traps, for the explanation of crystal phosphorescence.

trons can be bound at least temporarily. In the normal case, the *D* levels are empty. The particular position of the levels *C* and *D* with respect to the bands *A* and *B* depends on the ionization energy or electron affinity of the atoms under consideration. This explains why certain activator atoms are more favorable than others.

The mechanism of phosphorescence now seems to be the following: An electron is raised from the filled band *A* to the unfilled band *B* by the absorption of sufficiently shortwave light (or impact by  $\alpha$ -particles or electrons). In this state it seems to move freely in the lattice. In most crystals this absorption lies in the ultraviolet, since the separation of the energy bands is approximately equal to the first excitation step of the atom. This absorption by the fundamental lattice is very strong because *all* electrons in band *A* can absorb. The holes resulting from the ionization of electrons in band *A* are filled by a nonradiative transition from the occupied activator state *C*. The details of this transition have not yet been explained. The free electron in *B* now can diffuse to the location of an activator atom and thereby give up energy to the lattice, until it makes a transition to the state *C* under emission of light. On the other hand, it may, before it meets an activator atom, be trapped in a state *D*. We speak of trapping, for the transition  $D \rightarrow A$  is not

possible because the gaps in  $A$  are filled from  $C$ , and the transition  $D \rightarrow C$  is not possible because the separation of the two places in general is too great. Thus the electron and its excitation energy are stored in  $D$  until, by transfer of thermal energy or by the absorption of longwave radiation, it gets back to state  $B$ , diffuses to an activator atom, and emits light by a transition to  $C$ . This seems to be the essential mechanism of phosphorescence which thus clearly depends on the number of electron traps (states  $D$ ), the number of emitting atoms (states  $C$ ), and the temperature. The very longwave absorption  $D \rightarrow B$  is very weak because of the relatively few occupied  $D$  states; by measuring the corresponding absorption, the number of occupied  $D$  states can be determined. The trapping states  $D$ , which are necessary for phosphorescence, can also be reached by direct absorption from the fundamental band  $A$ . This absorption is very much weaker than the fundamental lattice absorption  $A \rightarrow B$ , but stronger than the longwave absorption  $D \rightarrow B$ , which depends on the number of occupied  $D$  states. There is, finally, the possibility of absorption by the activator atoms themselves (transition  $C \rightarrow B$ ). This absorption is proportional to the number of activator atoms ( $C$  states).

It is clear that the diffusion of the liberated electrons in band  $B$  represents an electric convection current and therefore, if an electric field is applied to the crystal, it can be measured electrically just as in the corresponding cases of the  $P$  electrons (page 459) or the photoelectrons in semiconductors.

No doubt, the theoretical picture described here is much too simple. It will require considerable development and refinement before it can explain in detail the mass of existing empirical results which is continuously increasing at a surprising rate.

Recently, for instance, phosphors with two kinds of activator atoms (e.g.,  $\text{SrSe}$  with  $\text{Cu}$  and  $\text{Sm}$ ) were investigated in more detail. While the copper atom as the "dominant activator" is responsible for the emission process, the auxiliary activator (in this case samarium) supplied the trapping states ( $D$  in Fig. 226). Leverenz, furthermore, distinguishes between  $s$  and  $i$  phosphors, according to whether the activator occupies an interstitial place or, as a substitutional atom, a normal lattice place. Many known differences in the behavior of different phosphors, for instance with respect to their photoconductivity, can be understood on the basis of this distinction.

We mention in concluding, that the macrocrystalline structure of phosphors seems to be responsible for their different behavior with respect to excitation by light or particles. High light efficiency from corpuscular excitation is obtained only with larger, reasonably perfect crystals in which the free electrons travel a sufficiently long mean free



path in the lattice until they find an activator atom. *Light* absorption with subsequent emission, on the other hand, occurs very frequently either by the activator atoms themselves (transition  $C \rightarrow B$ ) or in the close neighborhood of an activator atom, so that a good phosphorescence yield with light excitation is found in crystal powders in which the electrons cannot migrate very far.

### 7-16. Collision Processes at the Surface of Solids

To conclude our atomistic treatment of solid state physics, we shall briefly consider a number of interesting collision processes which occur at the surface of solids. First, there is the question of the possible ionization of atoms at the surface of an incandescent metal. Second, there are the phenomena connected with an energy transfer to the crystal lattice or crystal electrons in collisions between electrons, ions, normal or excited atoms on the one side, and the surface of a solid body on the other.

In discussing the ionization at an incandescent surface we start from the experimental results according to which alkali atoms in the vapor phase that impinge on an incandescent tungsten strip leave the strip as ions, if they are atoms of the heavy alkalis potassium, rubidium, and cesium. However, atoms of the lighter alkalis lithium and sodium leave the incandescent tungsten surface unionized. On the other hand, *all* alkali atoms are ionized at an incandescent platinum sheet. We may assume that the atoms which come in contact with the metal surface are at first adsorbed and then, as a result of the high temperature of the sheet, evaporate. Then it is only a question of whether a departing atom carries its outer electron with it or leaves it in the metal.

We can understand the whole situation if we imagine that the evaporation may occur in several steps (analogous to the Born cycle, page 429). Instead of the neutral atom we can imagine that first an ion evaporates. By adding energy equal to the work function  $\Phi$  we then release an electron from the metal and let it recombine with the ion to form an atom, thereby releasing the ionization energy  $E_i$ . Actually the process occurs in such a manner that the least possible energy is required.

Thus, if

$$\Phi > E_i \quad (7-19)$$

the electron remains in the metal and the atom leaves the surface as an ion. For in order to liberate an electron from the metal and let it recombine with the ion, the energy  $\Phi - E_i$  would have to be expended. In the normal case, however, where

$$E_i > \Phi \quad (7-20)$$

the energy  $E_i - \Phi$  is gained if an electron is "pulled" out of the metal and united with the ion. In this case the atoms will *not* be ionized in collisions with the incandescent metal. These theoretical conclusions are in good agreement with the experimental results. The heavy alkali atoms K, Rb, and Cs have ionization potentials which are smaller than the work function of tungsten, 4.5 ev. The light alkalis Li and Na, on the other hand, have larger ionization potentials. However, the ionization potentials of *all* alkalis lie below the work function of platinum with 6.2 ev; thus *all* alkali atoms are ionized by incandescent platinum.

The reverse process occurs at the cathode of glow and arc discharges. Following the lines of force of the electric field, positive gas ions from the discharge arrive at the cathode and are neutralized. This implies that they pull an electron out of the cathode and leave it as an atom. According to our previous discussion, this neutralization is energetically possible only if (7-20) is fulfilled. This, however, is the case, if we disregard discharges in the heaviest alkali vapors and a few electrode metals with particularly high work functions.

Now we shall discuss the mechanism of the energy transfer to the metal. For in our example the energy  $E_i$  is released in the neutralization of the arriving ion and is used partly for the release of a metal electron (under expenditure of the work function  $\Phi$ ), whereas the rest,  $E_i - \Phi$ , heats the cathode. This mechanism of neutralization may be regarded in a very pictorial manner. The positive ion attracts electrostatically a quasi-free metal electron and recombines with it. The process is more difficult to understand, if the condition

$$E_i > 2\Phi \quad (7-21)$$

is fulfilled. In this case the potential energy released in the neutralization of the ion is larger than the work function of *two* electrons. In these cases, actually two electrons are released. This secondary electron emission by ion bombardment is observed particularly in discharges in the noble gases helium and neon with their high ionization potentials. However, a metastable helium atom (see page 139) can also give up its potential energy (excitation energy) in a collision of the second kind with a metal surface and release one or two secondary electrons. The mechanism of this process is difficult to understand because no electrostatic effect is possible now, and the purely mechanical effect of the collision does not explain the emergence of an electron (or even two), because the impact is directed toward the interior of the metal. Basically the same difficulty is met in the transfer of purely kinetic energy to the metal in collisions. Fast electrons of sufficient kinetic energy, in collisions with metals, can release a large number of secondary electrons,

a technically very important process. Fast positive ions, similarly, can liberate, though with much smaller efficiency, electrons at the cathode of glow discharges, and furthermore cause the phenomenon of cathode sputtering, i.e., remove whole atoms from the metal surface.

It seems very probable that thermal processes and not simply mechanical collision effects are responsible for all these phenomena. We know, for example, that the anode of a high-current carbon arc and of certain spark discharges can be heated to violent evaporation by the discharge electrons which release on the anode their kinetic energy as well as the energy corresponding to the work function. The possibility has been considered that cathode sputtering is due to a similar effect of positive ions. However, this explanation seemed to be ruled out by the observation that the heat conductivity seems to have no influence on cathode sputtering. To cite one example: metallic silver and the poorly conducting silver chloride both exhibit equally strong cathode sputtering. However, all these difficulties are removed if, according to Seeliger, these phenomena are regarded as a very local overheating or, to use a familiar expression, as an extreme case of local lattice loosening. We have to imagine, then, that the temperature of a very small region of the cathode exceeds by far the critical temperature so that we have a gas-like state of matter of high density. Whether there exists an actual thermal equilibrium between the original lattice particles and the colliding particles seems more than doubtful, however.

In concluding this chapter on atomistic solid state physics, we again stress the fact that a satisfactory and rigorous presentation of this field, according to the present state of our knowledge, is not yet possible. However, we hope that our discussion of the most important problems of solid state physics has, on the one side, shown how fruitful the atomistic method of treatment has proved to be and that it may, on the other side, have helped us to realize how important it is to pay close attention in the future to this promising field of atomic physics.

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# APPENDIX

## *The Most Important Constants and Energy Relations of Atomic Physics*

Planck's constant.....	$h = (6.623 \pm 0.001) \times 10^{-27}$ erg-sec
Velocity of light.....	$c = (2.99776 \pm 0.00004) \times 10^{10}$ cm sec <sup>-1</sup>
Charge of the electron...	$e = (4.8024 \pm 0.0005) \times 10^{-10}$ esu $= (1.60199 \pm 0.00016) \times 10^{-19}$ coulombs (amp-sec)
Rest mass of the electron.	$m_0 = (9.1055 \pm 0.0012) \times 10^{-28}$ gm
Rest mass of the proton..	$M_p = (1.6723 \pm 0.0001) \times 10^{-24}$ gm
Rest mass of the neutron..	$M_N = (1.6747 \pm 0.0002) \times 10^{-24}$ gm
Atomic weight of the proton.....	$A_p = 1.007582 \pm 0.000004$
Atomic weight of the neutron in physical units..	$A_N = 1.008939 \pm 0.000007$
Atomic weight of the electron.....	$A_e = (5.4862 \pm 0.0006) \times 10^{-4}$
Avogadro's number.....	$N_0 = (6.0235 \pm 0.0004) \times 10^{23}$
Faraday.....	$F = 96522 \pm 7$ coulombs/mole
Boltzmann's constant....	$k = (1.3803 \pm 0.0001) \times 10^{-16}$ erg deg <sup>-1</sup>

Relations between different energy units used in atomic physics:

$$1 \text{ ev} = 8,067.5 \text{ cm}^{-1} = 1.602 \times 10^{-12} \text{ erg} = 23.04 \text{ kcal/mole}$$

$$1 \text{ atomic mass unit} = 931 \text{ Me}$$

Relation between excitation potential  $E$  (electron volts) and wavelength  $\lambda$  in angstroms of corresponding radiation:

$$\lambda(\text{\AA}) \times E(\text{ev}) = 12,400$$





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